

# Low Scale Flavor Gauge Symmetries

Benjamín Grinstein<sup>1,2\*</sup>, Michele Redi<sup>1†</sup> and Giovanni Villadoro<sup>1‡</sup>

<sup>1</sup>*CERN, PH-TH, CH-1211, Geneva 23, Switzerland*

<sup>2</sup>*Physics Dept., UCSD, La Jolla, CA 92093, USA*

## Abstract

We study the possibility of gauging the Standard Model flavor group. Anomaly cancellation leads to the addition of fermions whose mass is inversely proportional to the known fermion masses. In this case all flavor violating effects turn out to be controlled roughly by the Standard Model Yukawa, suppressing transitions for the light generations. Due to the inverted hierarchy the scale of new gauge flavor bosons could be as low as the electroweak scale without violating any existing bound but accessible at the Tevatron and the LHC. The mechanism of flavor protection potentially provides an alternative to Minimal Flavor Violation, with flavor violating effects suppressed by hierarchy of scales rather than couplings.

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\*bgrinstein@ucsd.edu

†michele.redi@cern.ch

‡giovanni.villadoro@cern.ch

# 1 Introduction

In the Standard Model (SM) the only source of flavor violation arises from the Yukawa couplings. In the limit of vanishing quark masses the SM Lagrangian acquires a large global symmetry (known as flavor or horizontal symmetry) mixing SM fermions of different generations. It is quite tempting to impose such a symmetry on the physics beyond the SM in order to suppress extra flavor violation. In the extreme case where all new physics effects are flavor universal up to small corrections satisfying the full flavor symmetry and proportional to the SM Yukawa, the scale of such new physics can be a TeV. This idea, known as minimal flavor violation (MFV) is quite old [1, 2], but recently is gaining ever more interest (see [3, 4]) as a consequence of the persistent failure to find flavor violation beyond the Standard Model.

The idea of assuming horizontal symmetries to be true symmetries of nature is even older [5]. Unfortunately such assumption itself is not enough to suppress flavor violation below the experimental bounds when the flavor symmetry is broken at low scales. The classical arguments against low-scale flavor symmetry work as follows (see, *e.g.*, [6] for a nice review). In order to produce the SM Yukawa couplings flavor symmetry must eventually be broken by the vacuum expectation value (VEV) of some field (“flavon”). This implies the presence of massless Goldstone bosons (GB) and bounds from hadron decays and astrophysics on such states are even stronger than those from flavor physics. Of course GBs can be easily avoided by gauging. In this case, however, there are flavor gauge bosons that mediate dangerous flavor changing neutral currents (FCNC) and their masses must be well above the TeV scale.<sup>1</sup>

Even requiring that the only sources of flavor breaking are the SM Yukawas is not enough to avoid large FCNC from the flavor gauge bosons. Indeed if the masses of the gauge bosons are proportional to the SM Yukawa couplings, they generate tree-level four-fermion operators proportional to inverse powers of the SM Yukawa couplings, enhancing FCNC among the first generations. Therefore despite the fact that the only spurions breaking flavor are the SM Yukawa couplings, as in MFV, inverse powers of the spurions appear in the higher dimensional operators, producing *de facto* “maximal” flavor violating operators. This argument shows how MFV models cannot arise directly from having a fundamental flavor symmetry in the underlying theory but rather from accidental ones.

In the classical argument against low-scale flavor gauge bosons sketched above we can easily identify a way out. If the fields breaking the flavor symmetry are instead proportional to the inverse of the SM Yukawa couplings, the effective operators generated by integrating out the flavor gauge bosons will be roughly proportional to positive powers of the Yukawa couplings, suppressing flavor violating effects for the light generations, much like in MFV models. The spectrum of the extra flavor states, controlled by the flavon VEVs, will thus present an inverted hierarchy, with states associated to the third generation much lighter than those associated to the first two. Models implementing this inverted hierarchy were first introduced in [9, 10].

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<sup>1</sup>Gauging only an abelian subgroup would not help because FCNC are reintroduced after the rotation needed to go to the mass eigenstate basis.

Remarkably, as we will show in this paper, the mechanism described above is automatic in the minimal extension of the SM with gauged flavor symmetries. With just the SM fermion content the full SM plus flavor gauge group would be anomalous. Extra flavorful fermions have to be added to cancel these anomalies. Such fermions are also welcome as they can make the SM Yukawa terms arise from a renormalizable Lagrangian, now that the Yukawa couplings have been uplifted to dynamical “flavon” fields. The smallest set of fermions canceling all anomalies leads automatically to the inverted-hierarchy structure mentioned above. The quantum numbers of these extra fermions are indeed such that the mixing with the SM fermion is flavor diagonal while their masses are proportional to the flavon VEVs. The SM fermion masses arise via a see-saw like mechanism, after integrating out the extra fermions, and are thus proportional to the inverse of the flavon VEVs. All non-SM particle masses (fermions, vector bosons and flavon fields) are controlled by the flavon VEVs, thus they are roughly proportional to the inverse of the SM Yukawa. The resulting inverted hierarchy in the new physics sector protects the SM fermions from getting large flavor breaking effects even when the lightest new states lie at the electroweak scale.

There are a number of analogies with MFV models: new physics effects are controlled by the flavor group, we may have models where only one spurion for each SM Yukawa matrix breaks the flavor symmetry, and the flavor breaking effects follow the hierarchical structure of the Yukawa couplings. However, this kind of models are not MFV. Indeed, in these models there is a limit where all Yukawa couplings vanish but flavor breaking effects remain finite. Contrary to the naive intuition that flavor violating effects must be larger than in MFV, it is very easy to find values of the parameters that produce extra flavor non-universal states without incurring into a flavor problem. Moreover, unlike in MFV, these could be light, even below the electroweak scale. The tightest bounds on this kind of models do not come from flavor breaking observables but rather from electroweak precision tests (EWPT) and direct searches for new particles, opening the possibility for direct discoveries of flavor physics at the LHC.

The mechanism protecting from flavor violations is robust against deformations of the model, both when more flavon fields are considered and when the nature of the flavor subgroup that is actually gauged is changed. Of course the detailed structure of the flavor sector as well as the size of the flavor violations will depend on these modifications but the latter will continue to remain sufficiently small in most of the parameter space of the theory.

We should point out that the possibility of having non-universal gauge bosons (and other flavorful physics) at the TeV scale is well known in the literature. In composite Higgs models and similar extra-dimensional constructions, for example, there is the possibility of having extra flavorful gauge bosons. Unlike in our case however, these vector fields are not “the” flavor gauge bosons (*viz.* they are not the states eating the Goldstone bosons of the broken flavor symmetry), but rather gauge bosons in some non-trivial representation of the flavor group, getting mass splitting from the breaking of the flavor symmetry (which generically is explicit in these models). In this respect they are closer to realizing the MFV idea (see, *e.g.*, [7, 8]).

In the rest of the paper we give the details on how the mechanism works, we will discuss where the strongest bounds on the model come from and possible signatures at hadron colliders. For definiteness we will focus on the quark sector, gauging the full flavor group and considering mainly the minimal set of flavon fields, although the same mechanism can easily be applied to more general situations.

## 2 Inverted Hierarchies From Anomaly Cancellation

In the absence of Yukawas, focusing on the quark sector, the SM enjoys at the classical level the global symmetry

$$U(3)_{Q_L} \otimes U(3)_{U_R} \otimes U(3)_{D_R}, \quad (2.1)$$

where  $Q_L$ ,  $U_R$  and  $D_R$  transform as fundamentals.

We assume this to be an exact symmetry of nature. In order to allow Yukawa couplings the flavor symmetry should be broken spontaneously by the vacuum. This can be most simply realized by the VEVs of two bifundamentals flavon fields transforming as

$$\begin{aligned} Y_u &= (\bar{3}, 3, 1), \\ Y_d &= (\bar{3}, 1, 3). \end{aligned} \quad (2.2)$$

In general the VEVs of these fields, while related, should not be confused with the Yukawa matrices, as functions of  $Y_{u,d}$  may have equal transformation properties. Indeed this will be the crucial feature of our model. To avoid problematic flavor violating GBs, the symmetry should be gauged. Within the SM the gauging of the SM flavor symmetry (2.1) is anomalous due to cubic and mixed hypercharge anomalies. The simplest option to cancel the cubic non-abelian anomalies is to add two right-handed colored fermions in the fundamental of  $SU(3)_{Q_L}$ , one left handed fundamental of  $SU(3)_{U_R}$  and one left-handed fundamental of  $SU(3)_{D_R}$ . In this way the fermions are vector-like with respect to the flavor gauge group but remain chiral with respect to the SM gauge symmetry. The other possibility, with the two right-handed triplets in an  $SU(2)_L$  doublet is an uninteresting, non-chiral model. We are therefore led rather uniquely to the following model:

	$SU(3)_{Q_L}$	$SU(3)_{U_R}$	$SU(3)_{D_R}$	$SU(3)_c$	$SU(2)_L$	$U(1)_Y$
$Q_L$	3	1	1	3	2	1/6
$U_R$	1	3	1	3	1	2/3
$D_R$	1	1	3	3	1	-1/3
$\Psi_{uR}$	3	1	1	3	1	2/3
$\Psi_{dR}$	3	1	1	3	1	-1/3
$\Psi_u$	1	3	1	3	1	2/3
$\Psi_d$	1	1	3	3	1	-1/3
$Y_u$	$\bar{3}$	3	1	1	1	0
$Y_d$	$\bar{3}$	1	3	1	1	0
$H$	1	1	1	1	2	1/2

Remarkably, with the above matter content all the anomalies except  $U(1)_{Q_L} \times SU(2)_L^2$  and  $U(1)_{Q_L} \times U(1)_Y^2$  automatically cancel. When, as required by cancellation of SM anomalies, the leptons are introduced  $U(1)_{B-L}$  remains anomaly free, so that  $U(1)_{Q_L}$  could also be gauged by gauging the  $B-L$  combination. The VEVs of  $Y_u$  and  $Y_d$  break  $U(1)_{Q_L} \times U(1)_{U_R} \times U(1)_{D_R}$  to the diagonal  $U(1)$  and an additional scalar field must be introduced in order to break also  $U(1)_{B-L}$  spontaneously. From now on we will focus on the gauging of  $SU(3)^3 \times U(1)^2$  which is the largest symmetry group broken by the SM Yukawa, other gaugings will be considered later.

The most general renormalizable Lagrangian reads,

$$\begin{aligned} \mathcal{L} = & \mathcal{L}_{kin} - V(Y_u, Y_d, H) + \\ & (\lambda_u \bar{Q}_L \tilde{H} \Psi_{uR} + \lambda'_u \bar{\Psi}_u Y_u \Psi_{uR} + M_u \bar{\Psi}_u U_R + \\ & \lambda_d \bar{Q}_L H \Psi_{dR} + \lambda'_d \bar{\Psi}_d Y_d \Psi_{dR} + M_d \bar{\Psi}_d D_R + h.c.), \end{aligned} \quad (2.3)$$

where  $M_{u,d}$  are universal mass parameters and  $\lambda_{u,d}^{(\prime)}$  are universal coupling constants. By a rotation of  $\Psi_u$  and  $\Psi_{uR}$  these parameters can be chosen to be real. The kinetic terms are built from covariant derivatives, which in our conventions are given by

$$DQ_L = \partial Q_L + ig_Q A_Q Q_L + ig_3 A_c Q_L + ig W Q_L + ig' \frac{1}{6} B Q_L \quad (2.4)$$

and similarly for the other fields.

In general, the VEVs of  $Y_{u,d}$  break the flavor symmetry to baryon number.<sup>2</sup> By a flavor transformation we can take  $Y_d = \hat{Y}_d$  diagonal and  $Y_u = \hat{Y}_u V$  where  $V$  is a unitary matrix. Integrating out the heavy fermions generates Yukawa interactions for the SM fields. At leading order for  $Y_{u,d} \gg M_{u,d}$  one immediately finds that the Yukawa couplings of the SM are

$$\begin{aligned} y_u &= V^\dagger \frac{\lambda_u M_u}{\lambda'_u \hat{Y}_u}, \\ y_d &= \frac{\lambda_d M_d}{\lambda'_d \hat{Y}_d}. \end{aligned} \quad (2.5)$$

Importantly the masses of the SM fermions follow an inverted hierarchy controlled by the inverse of  $\hat{Y}_{u,d}$  (see also [9, 10] for related works implementing the inverted hierarchy mechanism with models where the chiral diagonal  $SU(3)$  flavor symmetry is gauged). On the other hand, the exotic fermions have a mass proportional to  $\hat{Y}_{u,d}$  so that the lightest partner is the one associated to the top quark. As we will see this kind of see-saw mechanism is a general feature of the model through which all flavor and electroweak precision bounds can be easily avoided. The unitary matrix  $V$  plays the role of the CKM matrix of the SM. The formulas above receive important corrections for the third family since in this case the condition  $Y_{u,d} \gg M_{u,d}$  is not satisfied, particularly for the top quark. As we will see in the next section once this is properly accounted for it modifies the SM couplings. This produces important corrections to precision observables, in particular to the electroweak oblique parameters and the  $Zb\bar{b}$  coupling, which impose the most stringent bounds on the model.

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<sup>2</sup>We use the same notation both for the fields  $Y_{u,d}$  and their VEVs, except when the meaning is not immediate from the context.

## 2.1 Vectors and Scalars

The VEVs of  $Y_{u,d}$  give also a mass to the flavor gauge bosons,

$$\begin{aligned}\mathcal{L}_{mass} &= \text{Tr}|g_U A_U Y_u - g_Q Y_u A_Q|^2 + \text{Tr}|g_D A_D Y_d - g_Q Y_d A_Q|^2 \\ &= \frac{1}{2} V_{Aa} (M_V^2)^{Aa, Bb} V_{Bb},\end{aligned}\quad (2.6)$$

where

$$V_{Aa} = \{A_{Qa}, A_{Ua}, A_{Da}\}, \quad A_Q = A_{Qa} \frac{\lambda^a}{2}, \quad A_U = A_{Ua} \frac{\lambda^a}{2}, \quad A_D = A_{Da} \frac{\lambda^a}{2}, \quad (2.7)$$

$\lambda^{a=1,\dots,8}$  are the Gell-Mann matrices and  $\lambda^9$  is proportional to the identity.

The flavor gauge bosons couple to the quark currents,

$$J^{\mu ij, A} = (g_Q \bar{Q}_L^i \gamma^\mu Q_L^j, g_U \bar{U}_R^i \gamma^\mu U_R^j, g_D \bar{D}_R^i \gamma^\mu D_R^j). \quad (2.8)$$

Integrating out the vector fields SM four-fermion operators are produced, which in the flavor basis read

$$- \frac{1}{8} (M_V^2)^{-1}_{Aa, Bb} \lambda_{ij}^a \lambda_{hk}^b J_\mu^{ij, A} J^{\mu hk, B}. \quad (2.9)$$

In order to get the four-fermion operators in the mass eigenstate basis a further rotation by the unitary matrix  $V$  is needed on the left-handed up-quarks.

The flavor gauge bosons mediate FCNC since their masses break all flavor symmetries. Naively this implies the masses of all the gauge bosons to be around  $10^5$  TeV or higher in order to comply with flavor bounds. This expectation is however completely incorrect in our model because the masses depend on the inverse Yukawas. Roughly speaking the gauge bosons associated with transitions between light generation are automatically much heavier than the ones associated with the third generation with a hierarchy determined by the inverse Yukawas. As a consequence FCNC, which roughly scale as

$$\sim \frac{1}{Y_{u,d}^2} (\bar{q} \gamma^\mu q)^2, \quad (2.10)$$

are highly suppressed for the light generations.

To better understand how this works let us consider for simplicity the case where only  $Y_u$  is present. Since  $Y_u$  can be taken to a diagonal form there are no flavor violating processes and the individual family numbers are not broken so the associated gauge bosons remain massless. The masses of the flavor gauge bosons can be computed analytically in this case. Assuming equal couplings for  $SU(3)_{Q_L}$  and  $SU(3)_{U_R}$  the mass terms can be written as follows,

$$\begin{aligned}\mathcal{L}_{mass} &= \frac{1}{2} g^2 |V_{ij}|^2 (\hat{Y}_u^i - \hat{Y}_u^j)^2 + \frac{1}{2} g^2 |A_{ij}|^2 (\hat{Y}_u^i + \hat{Y}_u^j)^2 \\ &\approx \frac{1}{2} g^2 |V_{ij}|^2 \left( \frac{\lambda_u M_u}{\lambda'_u} \right)^2 \left( \frac{1}{y_u^i} - \frac{1}{y_u^j} \right)^2 + \frac{1}{2} g^2 |A_{ij}|^2 \left( \frac{\lambda_u M_u}{\lambda'_u} \right)^2 \left( \frac{1}{y_u^i} + \frac{1}{y_u^j} \right)^2,\end{aligned}\quad (2.11)$$

where  $V$  and  $A$  are the combinations  $(A_Q + A_U)/\sqrt{2}$  and  $(A_Q - A_U)/\sqrt{2}$  respectively. From this it follows that 4-fermion operators with light quarks obtained integrating out heavy gauge bosons

are very suppressed. The same mechanism works once the effects of  $Y_d$  are included, where all the flavor symmetries are broken and FCNC are generated. As we will show in various examples, flavor constraints can be avoided generically even if the lightest gauge boson is below the TeV scale.

The scalar sector is more model dependent due to the unspecified scalar potential. We discuss the radial fluctuations in detail in appendix A. After flavor symmetry breaking there are 10 radial fields contained in  $Y_{u,d}$ , corresponding to fluctuations of quark masses and CKM angles. These modes couple to fermion bilinears and therefore generate at low energy four-fermion operators. In particular fluctuations of the masses give rise to flavor diagonal operators and fluctuations of the CKM matrix induce flavor changing processes. However the suppression due to the inverted hierarchy works in this sector as well. To get an intuition for why this is the case we focus again on the flavor preserving four-fermion operators induced by the mass fluctuations. Their coupling to quarks is given by (for values of the couplings  $\lambda_{u,d}$  and  $\lambda'_{u,d}$  of order one)

$$\sim \frac{M}{\hat{Y}^i + \delta Y^i} v \bar{q}_i q_i \approx \left(1 - y^i \frac{\delta Y^i}{M}\right) m_{q_i} \bar{q}_i q_i, \quad (2.12)$$

so that the couplings of the radial modes  $\delta Y_i$  are highly suppressed for the light generations. Since these modes unitarize the scattering of the massive gauge bosons we expect their masses to be naturally set by the VEVs ( $m_{\delta Y_i} \sim \hat{Y}_i$ ). In this case the coefficients of the four-fermion operators scale as

$$\sim \left(\frac{y^i m_{q_i}}{M \hat{Y}^i}\right)^2 \quad (2.13)$$

which is extremely suppressed for the light generations. Actually the highly suppressed couplings alone would be enough to suppress dangerous four-fermion operators even when the flavon fields are light.

## 2.2 Remarks

A few comments are in order. While in MFV in the limit of vanishing Yukawa couplings the full flavor symmetry is restored, in our model there exists a limit where all Yukawa couplings vanish but flavor-breaking contributions remain finite. This can be seen by sending  $M_{u,d} \rightarrow 0$  with all other parameters fixed. In this limit  $y_{u,d} \rightarrow 0$  (see Eq. (2.5)) while four-fermion operators, depending only on  $Y_{u,d}$  (see Eq. (2.10)), still break flavor.

In the model above we assumed for simplicity the existence of only two bifundamentals. Actually in this case it can be shown that there is no renormalizable potential that gives rise to the Yukawa pattern of the SM. One possibility is to introduce non-renormalizable potentials. As long as the cut-off suppressing higher-dimensional operators is larger than the largest flavon VEV, its effects can be treated as perturbations, without spoiling our mechanism. Alternatively the  $Y_{u,d}$  could be combinations of several fields transforming as bifundamentals under the flavor group,

$$Y_{u,d} = \sum_{i=1}^N a_{u,d}^i X_{u,d}^i. \quad (2.14)$$

We have checked that in this case models with renormalizable potentials can be build. The mechanism of inverted hierarchy is still at work, leaving the fermion sector as before, however the relation between the Yukawas and the gauge boson masses is not uniquely determined. Unless the VEVs of the different fields are correlated the flavor gauge bosons will be generically heavier than in the minimal case improving flavor bounds but limiting the possibility of having these states at the electroweak scale.

In this paper we focus on the flavor symmetries of the quark sector but it is straightforward to extend this analysis to leptons at least when right-handed neutrinos are included. In this case the SM flavor symmetry is  $U(3)^6$ . For leptons cancellation of cubic anomalies works similarly to the quark sector and requires the addition of fermions transforming as singlets of  $SU(2)_L$  and with hypercharge opposite to the SM. In this way one finds that the only anomalous flavor symmetry is  $U(1)_{B+L}$ . We leave the detailed investigation of the lepton sector to future work.

One could also consider the gauging of smaller subgroups of the SM flavor symmetry. Obviously cancellation of anomalies can be achieved with the same matter content considered here so that the mechanism of inverted hierarchy works as before. An interesting subgroup is the diagonal  $SU(3)$  subgroup where the SM left- and right-handed fermions transform as fundamentals and anti-fundamentals respectively.<sup>3</sup> In this case however the mass of the  $SU(3)$  flavor gauge bosons is necessarily increased. Another interesting example is the gauging of abelian subgroups as also in this case, due to the inverted hierarchy, large corrections to FCNC do not arise.

Concerning unification the addition of the new fields charged under color and hypercharge worsens the unification of gauge couplings in the SM. Moreover in the case of  $SO(10)$  unification the flavor symmetry is only  $SU(3)$ . The simplest way to cancel the flavor cubic anomaly is to add fermions in the anti-fundamental representation of the flavor symmetry and the **16** of  $SO(10)$  to leave the theory chiral. However these degrees of freedom are insufficient to generalize our model since only one Yukawa term can be written down. Also for  $SU(5)$  the inverted hierarchy structure cannot be obtained at least in the simplest constructions. It is unclear to us how this model could be embedded in unification.

### 3 Experimental Bounds

In this section we consider the experimental bounds arising from the exotic fermions. These are the most model independent limits on the model as they only depend on four parameters, which can be conveniently chosen as the ratios  $\lambda_u/y_t$ ,  $\lambda_d/y_b$ ,  $M_u/m_t$  and  $M_d/m_b$ . The bounds originate from mixing effects between SM and exotic fermions that contribute in particular to  $Z \rightarrow b_L \bar{b}_L$ , EW oblique parameters,  $b \rightarrow s\gamma$  and  $V_{tb}$  as well as from direct searches.

In the previous section we integrated out the exotic fermions to leading order assuming  $Y_{u,d} \gg M_{u,d}$ . This is in general insufficient for the third family, in particular for the top whose large Yukawa

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<sup>3</sup>The other choice where the left and right fermions are fundamentals is already anomaly free within the SM and has been considered in the past, see for example [11] and Refs. therein.



requires a large or maximal mixing between SM and exotic fields and for the bottom whose coupling to  $Z$  may receive observable corrections.

The fermion mass matrices can be easily diagonalized in general. We can first eliminate the matrix  $V$  from the Yukawa interactions in eq. (2.3) by a simultaneous rotation of  $u_L$  and  $\Psi_{uR}$ . The physical states (without renaming the fields) are then given by the orthogonal rotations of the left and right fields,

$$\begin{pmatrix} u_R^i \\ u_R^{\prime i} \end{pmatrix} = \begin{pmatrix} c_{uRi} & -s_{uRi} \\ s_{uRi} & c_{uRi} \end{pmatrix} \begin{pmatrix} U_R^i \\ \Psi_{uR}^i \end{pmatrix}, \quad \begin{pmatrix} u_L^i \\ u_L^{\prime i} \end{pmatrix} = \begin{pmatrix} c_{uLi} & -s_{uLi} \\ s_{uLi} & c_{uLi} \end{pmatrix} \begin{pmatrix} U_L^i \\ \Psi_u^i \end{pmatrix}, \quad (3.1)$$

and similarly for the down-quark sector. The masses of SM and heavy fermions are then given by,

$$\begin{aligned} (m_u, m_c, m_t) &= \lambda_u \frac{v}{\sqrt{2}} \left( \frac{s_{uR1}}{c_{uL1}}, \frac{s_{uR2}}{c_{uL2}}, \frac{s_{uL3}}{c_{uL3}} \right), \\ (m_{u'}, m_{c'}, m_{t'}) &= M_u \left( \frac{c_{uL1}}{s_{uR1}}, \frac{c_{uL2}}{s_{uR2}}, \frac{c_{uL3}}{s_{uL3}} \right). \end{aligned} \quad (3.2)$$

We find it useful to define the physical variables,

$$x_i \equiv \frac{M_u}{m_{u_i}}, \quad y_i \equiv \frac{\lambda_u v}{\sqrt{2} m_{u_i}}, \quad (3.3)$$

which satisfy the properties,

$$x_i = \frac{c_{uR}^i}{s_{uL}^i}, \quad y_i = \frac{c_{uL}^i}{s_{uR}^i}, \quad \frac{m_{u_i'}}{m_{u_i}} = x_i y_i, \quad \frac{\lambda'_u \hat{Y}_u^i}{m_{u_i}} = \sqrt{(x_i^2 - 1)(y_i^2 - 1)}. \quad (3.4)$$

From the above relations one can easily derive,

$$s_{uLi} = \sqrt{\frac{y_i^2 - 1}{x_i^2 y_i^2 - 1}} = \frac{\lambda'_u \hat{Y}_u^i m_{u_i}}{\sqrt{(M_u^2 - m_{u_i}^2)^2 + (\lambda'_u \hat{Y}_u^i M_u)^2}}, \quad (3.5)$$

$$s_{uRi} = \sqrt{\frac{x_i^2 - 1}{x_i^2 y_i^2 - 1}} = \frac{\lambda'_u \hat{Y}_u^i m_{u_i}}{\sqrt{(\frac{1}{2}(\lambda_u v)^2 - m_{u_i}^2)^2 + \frac{1}{2}(\lambda_u \lambda'_u \hat{Y}_u^i v)^2}}. \quad (3.6)$$

Note that the physical region of  $x_i$  and  $y_i$  corresponds to  $x_i, y_i \geq 1$  or  $x_i, y_i < 1$ . In the first case  $m_{u_i'} \geq m_{u_i}$  while  $m_{u_i'} \leq m_{u_i}$  in the second. In the limit  $y_3 \rightarrow 1$  ( $\lambda_u \rightarrow y_t = \sqrt{2} m_t / v$ ), corresponding to  $\hat{Y}_u^3 \rightarrow 0$ , the right handed top becomes  $\Psi_{uR}^3$  while  $U_R^3$  becomes the right handed top-prime.

For phenomenological purposes and to better understand the parametric dependence of the results the following approximate expressions will be useful too:

$$\begin{aligned} s_{uLi} &= \frac{\lambda_u \lambda'_u v \hat{Y}_u^i}{\sqrt{2}(M_u^2 + \lambda_u^2 \hat{Y}_u^{i2})}, \\ s_{uRi} &= \frac{M_u}{\sqrt{M_u^2 + \lambda_u^2 \hat{Y}_u^{i2}}} \left( 1 - \frac{\lambda_u^2 \lambda_u'^2 v^2 \hat{Y}_u^{i2}}{2(M_u^2 + \lambda_u^2 \hat{Y}_u^{i2})^2} \right), \end{aligned} \quad (3.7)$$

valid up to terms  $\mathcal{O}(v^3)$  in the expansion in the SM Higgs VEV with respect to the new scales ( $\hat{Y}_t$  and  $M_u$ ). Note that before electroweak symmetry breaking only the right-handed quarks mix with the exotic fermions  $\Psi_{u,dR}$  while the left-handed mixing is suppressed by the Higgs VEV.

The most important consequence of the mixing is that the quark couplings are modified relative to those in the SM. For example the charged current (that couples to  $g_2 W_\mu^+/\sqrt{2}$ ) in terms of mass eigenstates is

$$\bar{u}_L(c_{u_L} V c_{d_L}) \gamma^\mu d_L + \bar{u}_L(c_{u_L} V s_{d_L}) \gamma^\mu d'_L + \bar{u}'_L(s_{u_L} V c_{d_L}) \gamma^\mu d_L + \bar{u}'_L(s_{u_L} V s_{d_L}) \gamma^\mu d'_L. \quad (3.8)$$

We have used the shorthands  $c_{u_L} = \text{Diag}(c_{u_{L1}}, c_{u_{L2}}, c_{u_{L3}})$ , etc, and  $V$  is the unitary matrix introduced in (2.5). Effectively the CKM matrix now becomes

$$V_{CKM} = c_{u_L} \cdot V \cdot c_{d_L}. \quad (3.9)$$

Note that such a matrix is not unitary. However, as we will see shortly, all the  $s_{qLi}$  are exceedingly small except, possibly, for that of the top quark. Moreover, the  $6 \times 6$  matrix of couplings to the charge current *is* unitary, hence exhibiting a generalized GIM mechanism.

The couplings of quarks to the photon are not modified, since they are protected by gauge invariance. And since the right handed quarks only mix with singlets of equal charge their couplings to the  $Z$  (proportional to their electric charge) are not modified either. The coupling of left handed quarks to the  $Z$  is now through the current

$$\bar{u}_L(T_3^u c_{u_L}^2 - s_w^2 Q_u) \gamma^\mu u_L + \bar{u}_L(T_3^u c_{u_L} s_{u_L}) \gamma^\mu u'_L + \bar{u}'_L(T_3^u s_{u_L} c_{u_L}) \gamma^\mu u_L + \bar{u}'_L(T_3^u s_{u_L}^2 - s_w^2 Q_u) \gamma^\mu u'_L + (u \rightarrow d), \quad (3.10)$$

where  $Q_{u(d)} = 2/3(-1/3)$  and  $s_w$  is the sine of the weak mixing angle. Using Eq. (3.7) we see that  $\delta g_{b_L}/g_{b_L} \sim (m_b/M_d)^2$ . The couplings of quarks to the Higgs are also modified relative to those in the SM:

$$\frac{1}{\sqrt{2}} \lambda_u h [-\bar{t}_L c_{u_L} s_{u_R} t_R + \bar{t}_L c_{u_L} c_{u_R} t'_R - \bar{t}'_L s_{u_L} s_{u_R} t_R + \bar{t}'_L s_{u_L} c_{u_R} t'_R] + (u \rightarrow d) + \text{h.c.} \quad (3.11)$$

### 3.1 Bounds from the down sector

In Fig. 1 we present the allowed region of parameter space for the down sector. The main bounds arise from the modified  $Zb\bar{b}$  coupling and direct searches described below. The green region is allowed by all measurements at 95% CL while the yellow region is model dependent.

#### 3.1.1 $R_b$

According to Eq. (3.10)  $Z$ -couplings are not universal. The heavier the quark the larger the effect, so for  $Z$  decays the largest and most sensitive deviation from the SM predictions is in  $R_b$ , the branching fraction to  $b$  quarks. At tree level we find

$$\frac{\delta \Gamma_{Zb\bar{b}}}{\Gamma_{Zb\bar{b}}} = -s_{d_{L3}}^2 \frac{2 + 4s_w^2 Q_d}{1 + 4s_w^2 Q_d + 8s_w^4 Q_d^2} \approx -2.3 s_{d_{L3}}^2, \quad (3.12)$$

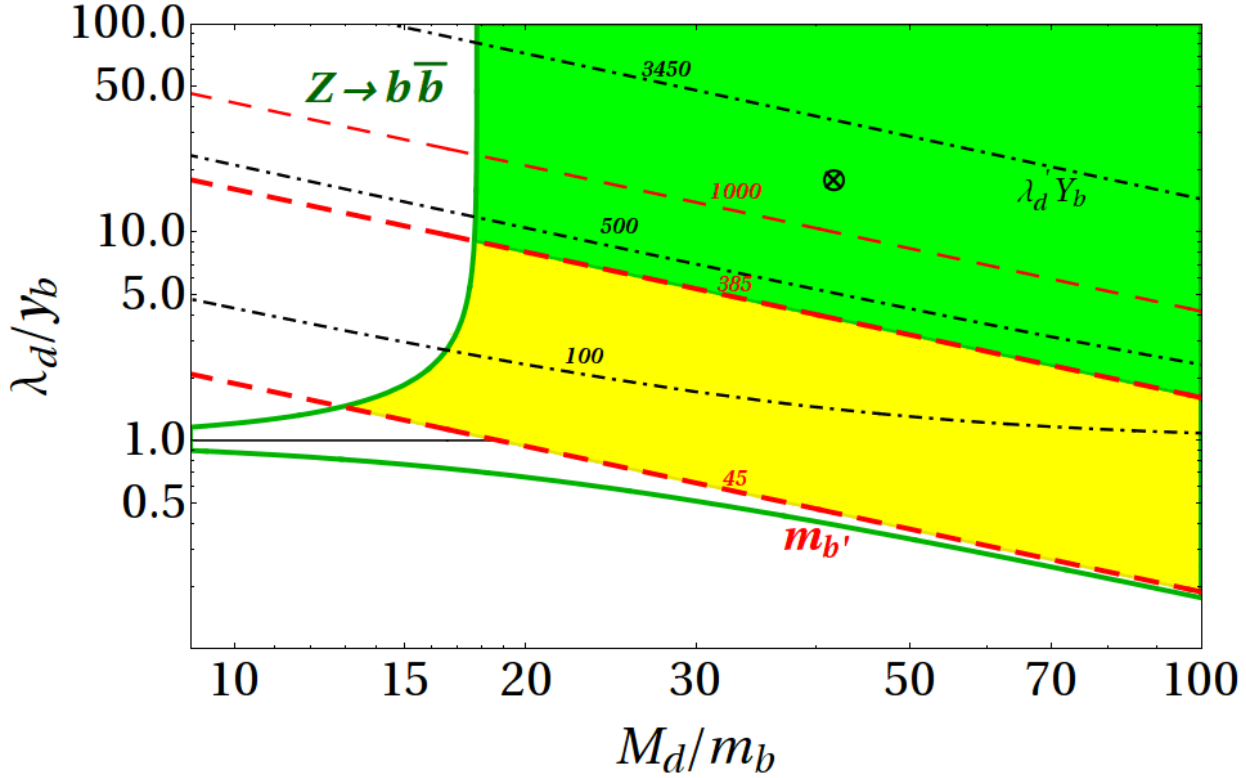


Figure 1: Allowed region of parameter space in the  $\lambda_d$  vs  $M_d$  plane. The yellow and green shaded regions are allowed by  $R_b$ , the thick green line labeled  $Z \rightarrow b\bar{b}$  corresponding to the 95% CL limit. The green one corresponds to  $m_{b'} > 385$  GeV, while the yellow one to  $45 \text{ GeV} < m_{b'} < 385$  GeV. Contours of constant  $m_{b'}$  (GeV) are shown in red dashed lines and contours of fixed  $\lambda'_d \hat{Y}_b$  (GeV) in black dash-dot lines. The black circle and cross show the choice of parameters in the examples of Sec. 5.

and writing  $\delta R_b/R_b^{\text{SM}} = (1 - R_b^{\text{SM}})(\delta\Gamma_{b\bar{b}}/\Gamma_{b\bar{b}}^{\text{SM}}) \approx 0.78(\delta\Gamma_{b\bar{b}}/\Gamma_{b\bar{b}}^{\text{SM}})$  we have

$$\frac{\delta R_b}{R_b^{\text{SM}}} \approx -1.8 s_{d_{L3}}^2, \quad (3.13)$$

to be compared to the current bound  $\delta R_b/R_b^{\text{SM}} \in [-4, 8] \cdot 10^{-3}$  at 95% CL [12].

Additional contributions to  $\delta R_b$  from couplings to light quarks are negligible. The virtual  $t$  and  $t'$  contributions deviate from the SM's virtual  $t$  contribution by an amount that vanishes both with  $m_{t'} - m_t$  and with  $s_{u_{L3}}^2$ . The resulting bound on these parameters is weaker than bounds presented below from  $V_{tb}$  (and the direct limit on  $m_{t'}$ ).

Fig. 1 shows the 95% CL bound from  $\delta R_b$  in the  $\lambda_d/y_b$  vs  $M_d/m_b$  plane, where  $y_b = \sqrt{2}m_b/v$ .

### 3.1.2 Direct bounds on $m_{b'}$

CDF data excludes a  $b'$  with mass above 100 GeV and below 268 GeV assuming  $\text{BR}(b' \rightarrow Zb) = 100\%$  [13, 15]. For masses above  $m_{b'} = m_t + M_W = 253$  GeV the  $Wt$  channel opens up and CDF data sets a mass limit  $m_{b'} > 385$  GeV assuming  $\text{BR}(b' \rightarrow Wt) = 100\%$  [16]. In our model the branching fraction assumptions may not apply. The couplings of the  $b'$  to  $Wt$  and  $Zb$  include a suppression factor of  $s_{d_{L3}} \lesssim 0.04$  (from  $R_b$ ). For a light Higgs the channel  $b' \rightarrow bh$  can become important. According to Eq. (3.11) the  $bb'$  couplings to the Higgs are  $\frac{1}{\sqrt{2}}\lambda_d c_{d_{L3}} c_{d_{R3}} = s_{d_{L3}} c_{d_{L3}} m_{b'}/v$  and  $\frac{1}{\sqrt{2}}\lambda_u s_{d_{L3}} s_{d_{R3}} = s_{d_{L3}} c_{d_{L3}} m_b/v$ . Hence  $\text{BR}(b' \rightarrow bh)$  will be large in the region where it is kinematically allowed, provided  $m_{b'} \gtrsim M_Z$ . The LEP2 95%CL bound on the Higgs mass,  $m_h > 114$  GeV, is valid in this model since the properties of a light Higgs are largely unchanged from that of the SM. Hence  $100 \text{ GeV} < m_{b'} \lesssim m_b + m_h \approx 118 \text{ GeV}$  is excluded. A bound  $m_{b'} > 128 \text{ GeV}$  is given by the PDG [12] based on D0 data [17] on  $WW + 2\text{jets}$ , used in top pair production searches. However, the bound assumes  $\text{BR}(b' \rightarrow Wq) = 100\%$ . The region between the LEP bound  $m_{b'} \gtrsim M_Z/2$  and  $m_{b'} < 100 \text{ GeV}$  is not easily excluded. D0 has excluded a 4th generation sequential charge  $-1/3$  quark up to  $m_{b'} = m_b + M_Z$  by searching for radiative decays  $b' \rightarrow b\gamma$  [18] (see also [19] for bounds using  $b' \rightarrow b\ell^+\ell^-$  from analysis of Tevatron data). These bounds again may not apply in our model, since the branching fractions are not those of a sequential fourth generation quark. For example, the tree level three body decay  $b' \rightarrow h^*b \rightarrow b\bar{b}b$  can compete well with the two body radiative decay. The yellow and green shaded regions in Fig. 1 are allowed by  $R_b$ , the thick green line labeled  $Z \rightarrow b\bar{b}$  corresponding to the 95% CL limit. The green one corresponds to  $m_{b'} > 385 \text{ GeV}$ , while the yellow one to  $45 \text{ GeV} < m_{b'} < 385 \text{ GeV}$ , and may or may not be excluded depending on the value of the Higgs mass and, to lesser extent, other model parameters, *e.g.*, flavon masses. For reference the figure shows contours of constant  $m_{b'}$  in red dashed lines and contours of fixed  $\lambda'_d \hat{Y}_b$  in black dash-dot lines.

## 3.2 Bounds from the up sector

Experimental bounds on the up sector are collected in Fig. 2. The physical region of parameters corresponds to the first and third quadrants where  $m_{t'} \geq m_t$  and  $m_{t'} \leq m_t$ , respectively. The main constraint in the first region arises from precision electroweak constraints and in particular from corrections to the  $T$  parameter. The second region (where constraints from  $T$ ,  $S$  and  $U$  are not applicable) is strongly constrained by  $V_{tb}$ ,  $b \rightarrow s\gamma$  and direct searches.

### 3.2.1 Electroweak Precision Tests

The exotic fermions modify the oblique corrections to the electroweak gauge bosons with respect to their SM values. We compute the oblique parameters  $S$ ,  $T$  and  $U$  in appendix B. Since the exotic fermions are  $SU(2)_L$  singlets they only contribute through the mixing with SM left doublets.

After electroweak symmetry breaking the quark doublets mix with the left singlets. This violation

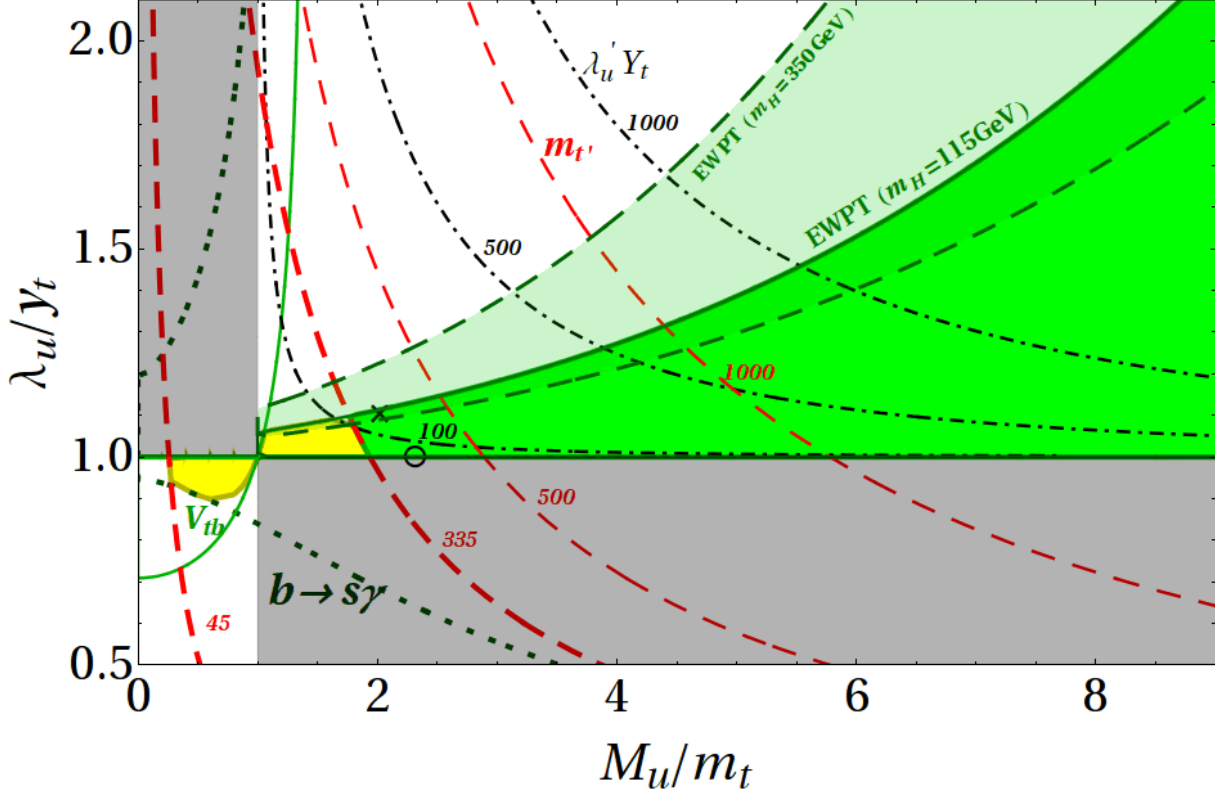


Figure 2: Allowed region of parameter space in the  $\lambda_u$  vs  $M_u$  plane. The shaded grey region is unphysical. The thick green line labeled EWPT shows the region allowed at 95% CL by the EW oblique parameters for  $m_H = 115$  GeV. For  $m_H = 350$  GeV the allowed region becomes the one between the green dashed lines. The thin green line labeled  $V_{tb}$  shows the 95% CL limit from direct single top production while the green short-dashed line shows the 95% CL bound from  $b \rightarrow s\gamma$ . Of the region allowed by EWPT,  $V_{tb}$  and  $b \rightarrow s\gamma$  we have distinguished  $m_{t'} > 335$  GeV shaded in green from  $45 \text{ GeV} < m_{t'} < 335 \text{ GeV}$ , shaded in yellow. For the latter direct mass bounds may (or not) apply, depending on the Higgs mass and other model parameters. Contours of constant  $m_{t'}(\text{GeV})$  in red dashed lines and contours of fixed  $\lambda'_u \hat{Y}_t(\text{GeV})$  in black dash-dot lines. The black circle and cross show the choice of parameters in the examples of Sec. 5.

of custodial symmetry generates a correction to the  $T$ -parameter.<sup>4</sup> For simplicity we only consider the contributions of the third family, which are the dominant ones. In the limit  $m_b \rightarrow 0$  the exact one

<sup>4</sup>This was also studied recently in [14] in a model with vector like top partners. For the third generation fermions our model reduces to theirs.

loop formula derived in App. B reads

$$T = \frac{3 s_{uL3}^2}{8\pi s_w^2 c_w^2} \frac{m_t^2}{M_Z^2} \left[ c_{uL3}^2 \left( \frac{m_{t'}^2}{m_t^2 - m_t^2} \log\left(\frac{m_{t'}^2}{m_t^2}\right) - 1 \right) + \frac{s_{uL3}^2}{2} \left( \frac{m_{t'}^2}{m_t^2} - 1 \right) \right]. \quad (3.14)$$

As explained above for  $s_{uL3} = 0$  the correction to  $T$  vanishes. From Eq. (3.3) this corresponds to  $M_u \rightarrow \infty$  or  $\lambda_u = y_t$  (*i.e.*,  $\hat{Y}_u^3 = 0$ ). In the first case the exotic fermions acquire an infinite vector like mass so the correction to  $T$  obviously vanishes in this limit. In the second case the mass of the top partner can be light. Since the amount of custodial symmetry breaking is proportional to  $\lambda_u$  we expect  $T$  to have the same sign as  $\lambda_u - y_t$ , as is readily checked using the explicit formulas. Even though the contribution to  $T$  is smaller than in a fourth generation model, it can be sizable and gives one of the most important bound on the parameters of the model.

On the other hand, the contribution of the exotic fermions to the  $S$ -parameter is always small and its sign is not fixed. This is similar to four-generation models, where corrections to  $S$  are generically smaller than to  $T$ . Despite the small contribution to  $S$  the bound obtained combining  $S$  and  $T$  is significantly more restrictive than the one from  $T$  alone, due to the correlation between  $S$  and  $T$  in the electroweak fit. In the allowed region in Fig. 2 we have also included the bound from the  $U$  parameter which however only affects the results in a minor way.

A few words of caution. The new physics contributions to precision electroweak parameters are here obtained as the difference between our model and the SM one loop value. The mixing also modifies the two loop SM correction which is not negligible in the SM. This effect is relevant in the region where  $s_{uL3}$  is large which is however only allowed in the region of small  $m_{t'}$  ( $\ll$  TeV). In this region, however, the canonical  $S, T, U$  parameters are in general insufficient and a more refined analysis is needed. Moreover our bounds are obtained with the assumption of a light Higgs. In the SM, increasing the Higgs mass would worsen the global electroweak fit mainly because of the negative contribution to the  $T$  parameter (the contribution to  $S$  is instead smaller and positive, and thus well within the bound). Interestingly in our model the correction to  $S$  is always small while the contribution to  $T$  is positive and easily of the right order to accommodate also an heavy Higgs. In Fig. 2 we have also shown the region of parameter space allowed for a Higgs with  $m_H = 350$  GeV which requires a non zero mixing. Therefore the bounds from oblique parameters should be taken with a grain of salt since both the Higgs mass and other new physics not related to flavor may alter the bounds.

### 3.2.2 $V_{tb}$

The effective CKM matrix in Eq. (3.9) is not unitary. Unitarity of the CKM matrix is presently only tested with significant accuracy on the first row,  $\sum_{q=d,s,b} |V_{uq}^{\text{CKM}}|^2 = 0.9999 \pm 0.0011$  [12]. However, since only light quarks participate in this, the resulting bound is very weak,  $M_{u,d}$  greater than a few GeV.

Unitarity of the third row is more restrictive. The measured smallness of  $|V_{td}^{\text{CKM}}|$  and  $|V_{ts}^{\text{CKM}}|$  implies that the unitary matrix  $V$  in Eq. (3.9) has  $|V_{tb}| = 1$  to high accuracy. Hence direct measurement

of the  $tbW$  coupling constrains  $c_{uL3}$ . Single top production experiments at the Tevatron set a 95%CL bound  $|V_{tb}^{\text{CKM}}| \approx c_{uL3} > 0.77$  [20]. The resulting constraint on the model parameters is shown in Fig. 2. The allowed values for  $c_{uL3}$  at 95% CL lie between the green line labeled  $V_{tb}$  and the one at  $\lambda_u/y_t = 1$ .

### 3.2.3 $b \rightarrow s\gamma$

There are two distinct underlying processes that give rise to radiative  $B$  decays. On the one hand there are  $\Delta B = -\Delta S = \pm 1$  operators, such as  $(\bar{s}b)(\bar{b}b)$ , produced by the exchange of either a flavor vector meson or a radial mode associated to the flavons. These contributions are highly model dependent and also very suppressed by the overall coefficient of the four quark operator.

On the other hand there are sizable and less model dependent contributions from SM-like graphs with a virtual  $Wt$  or  $Wt'$ . For  $m_{t'} > m_t$ , their sum is always larger than that of the SM. To see this note that if the SM amplitude is a function  $f(m_t)$  then the corresponding amplitude in this model is  $c_{uL3}^2 f(m_t) + s_{uL3}^2 f(m_{t'}) = f(m_t) + s_{uL3}^2 (f(m_{t'}) - f(m_t))$ . Since  $f(m_t)$  is a monotonically increasing function, the deviation from the SM result,  $s_{uL3}^2 (f(m_{t'}) - f(m_t))$  has the same sign as  $m_{t'} - m_t$ .

In more detail, working at NLO in the simplified but accurate approximation of Ref. [21] we find for the process  $b \rightarrow s\gamma$  that

$$\frac{\delta\Gamma_{bs\gamma}}{\Gamma_{bs\gamma}} = 2s_{uL3}^2 \frac{A(x') - A(x) - \frac{8}{3}(1 - z^{\frac{2}{23}})(D(x') - D(x))}{A(x) - \frac{8}{3}(1 - z^{\frac{2}{23}})D(x) - \frac{6}{19}X_2(1 - z^{\frac{19}{23}})}, \quad (3.15)$$

where  $x = m_t^2/M_W^2$  and  $x' = m_{t'}^2/M_W^2$  are arguments of the loop functions  $A$  and  $D$  (given in Ref. [21])  $z = \alpha_s(m_b)/\alpha_s(M_W)$  and  $X_2 = 232/81$  is the coefficient of anomalous dimension mixing the four quark operator into the transition magnetic moment operator. The resulting 95% CL bound in the  $\lambda_u/y_t$  vs  $M_u/m_t$  plane, where  $y_t = \sqrt{2}m_t/v$  is shown as a green short-dashes line in Fig. 2.

### 3.2.4 Bounds from $m_{t'}$

Fig. 2 also shows, as red dashes, contours of fixed  $m_{t'}$ . CDF excludes  $m_{t'} < 335$  GeV at 95%CL, assuming  $\text{BR}(Wq) = 100\%$  [22]. As discussed above for the case of the  $b'$  the branching fraction assumptions in the experimental analysis may not apply in this model. Of the region allowed by EWPT,  $V_{tb}$  and  $b \rightarrow s\gamma$  we have therefore distinguished  $m_{t'} > 335$  GeV shaded in green from  $45 \text{ GeV} < m_{t'} < 335$  GeV, shaded in yellow. For the latter direct mass bounds may (or not) apply, depending on the Higgs mass and other model parameters. For reference the figure shows contours of constant  $m_{t'}$  in red dashed lines and contours of fixed  $\lambda'_u \hat{Y}_t$  in black dash-dot lines.

## 3.3 Neutron EDM

The interactions among quarks due to flavor-vector or flavon exchange can give contributions to the Electric Dipole Moment (EDM) of hadrons. In the SM the dominant mechanism for EDM of the

neutron is from a  $\Delta S = -1$  CP violating transition  $n \rightarrow \Lambda, \Sigma^0$  followed by a  $\Delta S = 1$  transition  $\Lambda, \Sigma^0 \rightarrow n\gamma$  [23]. The CP violating interaction is a 1-loop induced four-quark “penguin” operator

$$\mathcal{L}_{\text{CPV}} = i \frac{3\alpha_s G_F}{9\sqrt{2}\pi} \ln(m_t^2/m_c^2) \text{Im}(V_{td}^* V_{ts}) (\bar{d}_L \gamma^\mu T^a s_L) \sum_{q=u,d,s} (\bar{q} \gamma_\mu T^a q). \quad (3.16)$$

A recent estimate gives [24]

$$d_n^{\text{SM}} \simeq 10^{-32} e \text{ cm}. \quad (3.17)$$

A rough estimate of the new contributions to the neutron EDM induced by flavor-vector or flavon exchange is obtained by replacing the coefficient of the four-quark “penguin” operator in the calculation of  $d_n^{\text{SM}}$  by the CP violating coefficient  $C_{\text{CPV}}$  of a newly induced four-quark operator. One then has

$$\Delta d_n \sim \frac{C_{\text{CPV}}}{\frac{3\alpha_s G_F}{9\sqrt{2}\pi} \ln(m_t^2/m_c^2) \text{Im}(V_{td}^* V_{ts})} d_n^{\text{SM}} \approx 7.6 \times 10^9 \text{ GeV}^2 C_{\text{CPV}} d_n^{\text{SM}}. \quad (3.18)$$

The resulting bounds on the model parameters are extremely weak. For example, the CP violating part of coefficients of  $\Delta S = \pm 1$  four-quark operators in Eq. (2.9) are numerically of order  $C_{\text{CPV}} \sim 10^{-15} \text{ GeV}^{-2}$ . The smallness of this result justifies the crude nature of the estimate (in which we have ignored, for example, the different possible Dirac and color structures that may arise in the four-quark operators). The coefficients of operators from flavon exchange, although more model dependent, are similarly small.

Additional contributions arise from graphs involving only electroweak interactions but in which the heavy quarks participate. These are all at best of the order of the SM contributions (for example, by modifying the coefficient of the penguin operator).

We conclude then that this class of models predicts small EDMs of hadrons, comparable in order of magnitude to those of the SM.

## 4 Signatures

Despite the small number of extra parameters beyond the SM ones and the relatively rigid structure of the spectrum of the model, mostly fixed by the SM Yukawa couplings, the phenomenology above the production threshold of new states is very rich, drastically changing in different regions of the parameters space. We will not attempt to cover here this subject, which deserves a separate study. Instead we will only give a sampling of some possible new signatures of the model (for recent more detailed analysis on similar models see, *e.g.*, [14, 25, 26], keeping in mind however that the BR in our case could be altered by the presence of the extra vector and flavon fields).

Among all new states, the one that presents less model dependence is the  $b'$ . As discussed before, existing searches do not provide very strong bounds on such particles and, as shown explicitly in the next section, it is easy to find parameters of the model where such resonance is within the reach of hadron colliders. The strong bounds from  $Z \rightarrow b\bar{b}$  force the  $b'$  to have small mixing with the SM  $b$



quark, which implies a small coupling with the  $SU(2)_L$  gauge fields. It turns out that, choosing  $O(1)$  values for the couplings of the model (such as  $\lambda_d$ ,  $\lambda'_d$  and the gauge couplings), the standard fourth-generation channels  $Wt$ ,  $Zb$  and  $Wt'$  can compete with others such as  $Z'b$ ,  $\tilde{b}b$  and  $bh$ . In particular the BR to  $bh$  can easily be of order one. Being a colored object, the  $b'$  could be pair produced copiously at the LHC, provided its mass is not too high. The signature would be quite striking having up to six bottom quarks in the final state ( $p\bar{p} \rightarrow \bar{b}'b' + X \rightarrow 2h + 2b + X \rightarrow 6b + X$ ).

For the  $t'$  the discussion is similar, with the possibility however of a substantial difference. In this case the bound on the mixing angle  $s_{u_{L3}}$  is weaker, coming only from EWPT. Choosing as before  $O(1)$  values for the parameters, we have two means of maintaining  $s_{u_{L3}}$  below the bounds, either by increasing  $M_u$  with respect to  $m_t$  or by suppressing  $\hat{Y}_u^3$ . In the first case we get a similar result to that of the  $b'$ , with the  $t' \rightarrow th \rightarrow W + 3b$  decay channel becoming important. The  $t'$  could be pair produced at hadron colliders, leading to very clean  $6b + WW$  signals (see [25] for a detailed study). Notice that both in the  $b'$  and in the  $t'$  case the  $O(1)$  BR into  $hb$  and  $ht$  could substantially increase the Higgs production cross section, improving the capability of discovering and studying its properties. In the second case, in the limit of small  $\hat{Y}_u^3$ , also the right mixing angle  $c_{u_{R3}}$  get suppressed. In this case the dominant channel becomes  $t\tilde{t}$ , if the radial mode of the top Yukawa ( $\tilde{t}$ ) is light enough. Actually in the limit  $\hat{Y}_u^3 \rightarrow 0$  the  $t'$  almost decouples from the SM, and an approximate discrete symmetry (similar to R-parity) prevents the lightest among the  $t'$  and the  $\tilde{t}$  from decaying into SM particles.<sup>5</sup> Such discrete symmetry is broken only by the mixing of the  $t'$  sector with the other generations, thus producing a highly suppressed decay rate for the  $t'$  in this scenario.

Finally some comments on the possible lightest gauge flavor fields. The lightest state is expected to couple more to the third generation, and in particular to the top. The actual couplings, however, depend on which flavor group is gauged and the magnitude of its coupling constants.

For the lightest states three possibilities are favored. If the gauge group is  $SU(3)^3$  then the lightest state couples through the diagonal Gell-Mann  $\lambda^8$  generator of  $SU(3)^3$ , thus with doubled strength to the third generation respect to the first two. In this case we get a leptophobic non-universal  $Z'$ , which can be produced directly via  $q\bar{q}$  annihilation and can decay either into  $t\bar{t}$  or into two jets. Existing Tevatron studies of similar  $Z'$  set mass bounds below a TeV [27]. More possibilities arise when the  $U(1)$ s are also gauged. In particular the flavor boson can mix with the SM hypercharge vector and acquire a coupling to leptons too. If such mixing is large, the lightest vector behave as a heavy  $Z'$  coupled to the hypercharge current and with an anomalous coupling to the right-handed top. In this case strong bound are present from the EWPT [28]. If instead the mixing with the hypercharge is negligible, then the lightest gauge boson will only couple to  $t'_L$  and to a linear combination of  $t_R$  and  $t'_R$  (depending on  $s_{u_{L3}}$ ). In this case the four-top(top') signal becomes one of the most interesting (see, *e.g.*, [29]).

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<sup>5</sup>Actually a similar limit is also possible for the down sector, when  $\hat{Y}_d^3 \rightarrow 0$ ,  $s_{d_{L3}} \rightarrow 0$  and the  $b'$  decouples from the  $b$ ; however this happen in the small coupling limit  $\lambda_d \rightarrow y_b \approx 1/40$ , then the scales of the  $s'$  and  $d'$  decrease accordingly and FCNC induced by the first generation may start becoming important.

## 5 Examples

The details of a particular realization of the mechanism described in Sec. 2 depend strongly on the actual model and parameters chosen. Depending on the gauge group ( $U(3)^3$ ,  $SU(3)^3 \times U(1)^2$ ,  $SU(3)^3$ ,  $SU(3)$ ,  $U(1)^n, \dots$ ), the number and representations of scalar flavon fields and the different parameters of the Lagrangian, the spectrum of the new particles and their couplings may vary substantially. Still there are some features that are rather model independent and characterize the model.

As shown before, with the exception of the top quark sector, the structure of the fermionic part of the model is quite rigid, depending only on the two scales  $M_u$  and  $M_d$ , the rest being fixed by the SM Yukawa couplings. Once the gauge group and the scalar content has been chosen so is the basic structure of the spin-1 sector. But as a result of the larger number of parameters connecting its spectrum and couplings to the SM Yukawa terms, such as the gauge couplings and extra Yukawa couplings ( $\lambda_{u,d}$ ,  $\lambda'_{u,d}$ ), it is far from being specified in detail.

In the following we will provide two explicit examples where all the parameters have been fixed, in order to demonstrate how easy it is to build explicit models with  $O(1)$  couplings, new flavor non-universal states at the TeV scale and compatibility with all existing experimental bounds. In fact, depending on the choice of the parameters the strongest bounds may come from different sources, such as EWPT,  $Z \rightarrow b\bar{b}$ , single top production at Tevatron,  $Z'$  searches and other direct bounds for spin-1 and spin-1/2 particles,  $\Delta M_K$ , etc...

The two examples below correspond to the two different flavor gaugings  $SU(3)^3$  and  $SU(3)^3 \times U(1)^2$ , respectively. For definiteness in both cases the flavon content have been chosen to be minimal: just the two  $Y_u$  and  $Y_d$  fields of Sec. 2. The couplings have been chosen to be  $O(1)$  and the two mass scales  $M_u$  and  $M_d$  to be low enough to produce interesting physics for high-energy colliders and possibly for next generation flavor experiments.

### 5.1 First example: An $SU(3)^3$ model

In the first example we choose the following parameters:

$M_u$ (GeV)	$M_d$ (GeV)	$\lambda_u$	$\lambda'_u$	$\lambda_d$	$\lambda'_d$	$g_Q$	$g_U$	$g_D$
400	100	1	0.5	0.25	0.3	0.4	0.3	0.5

Given the parameters above the entries of the flavon VEVs are fixed by requiring the right SM Yukawa couplings be reproduced, this gives<sup>6</sup>:

$$\begin{aligned} Y_u &\approx \text{Diag} (1 \cdot 10^5, 2 \cdot 10^2, 8 \cdot 10^{-2}) \cdot V \text{ TeV}, \\ Y_d &\approx \text{Diag} (5 \cdot 10^3, 3 \cdot 10^2, 6) \text{ TeV}, \end{aligned} \tag{5.1}$$

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<sup>6</sup>The values of the  $Y_{u,d}$  VEVs (and the the results that follow) have been calculated taking into account the running of the Yukawa couplings only up to the TeV scale. The effects coming from the running from the TeV scale up to the flavor breaking scales are more model dependent and affect mainly the value of the highest  $Y_{u,d}$  VEVs, which we do not need to know with high accuracy. In fact the knowledge of the order of magnitude for these quantities is enough for our purposes.

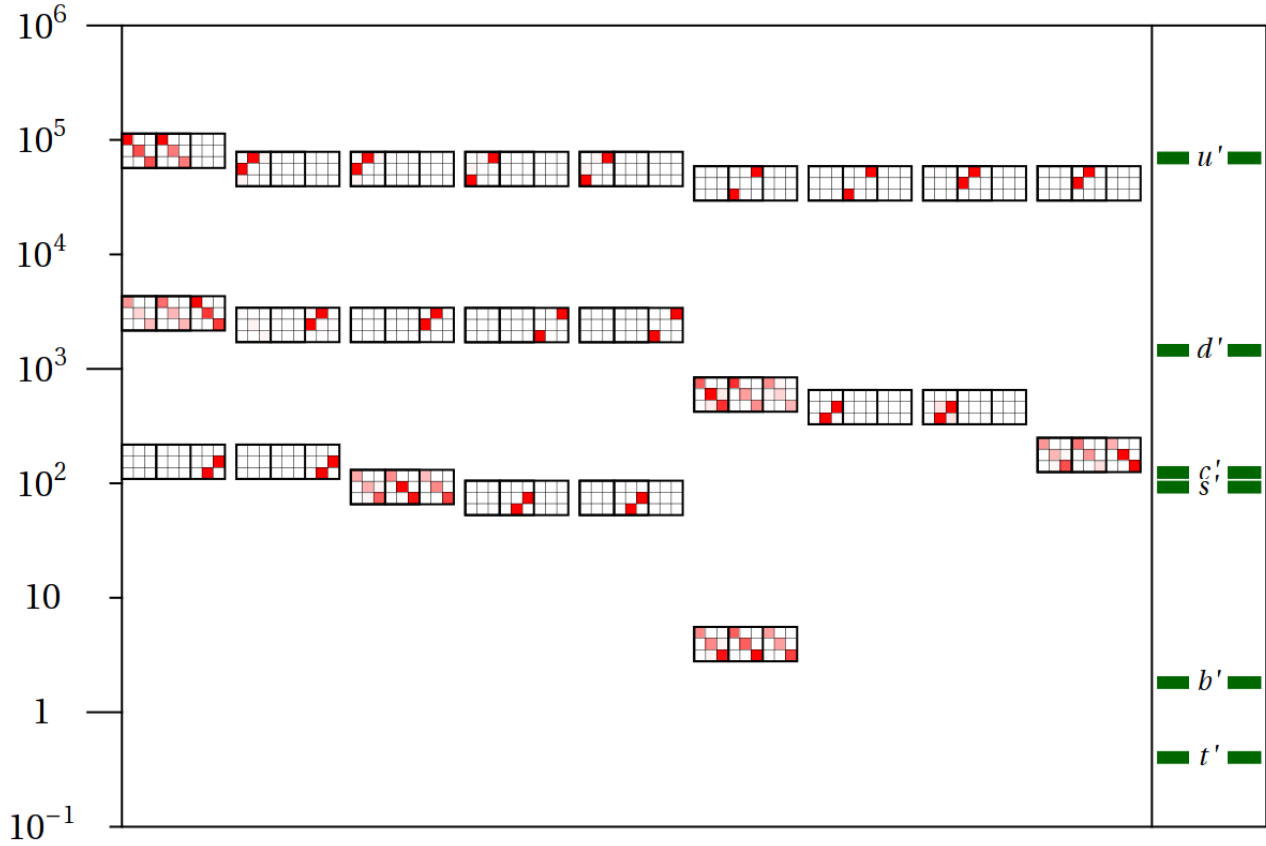


Figure 3: Spectrum of the flavor spin-1 (left) and spin-1/2 (right) fields for the first example (see text for details). Each vector fields is represented by a set of three  $3 \times 3$  matrices representing the associated generators to the three gauged  $SU(3)$  groups ( $SU(3)_Q$ ,  $SU(3)_U$ ,  $SU(3)_D$  respectively), the intensity of the color (from white to red) correspond to the size of each entry in the generators (from 0 to 1). The position in the vertical axis represent instead the corresponding mass in TeV, analogously for the masses of the heavy quark partners, on the right.

where  $V$  is the unitary experimental CKM matrix [12].

The couplings are chosen to be smaller than 1 to avoid possible problems with early Landau-poles except for  $\lambda_u$ , which must be larger than  $y_t = \sqrt{2}m_t/v \simeq 1$  (or slightly smaller when  $m_{t'} < m_t$ ; see Sec. 3). For  $\lambda_u = 1$ , as in this example, the mixing of the left doublet is small and the lowest eigenvalue of  $Y_u$  approaches zero.

Given the parameters above we can calculate both the spectrum and couplings of the spin-1 and spin-1/2 sectors of the theory. The spectrum is summarized in Fig. 3.

The masses of the four lightest spin-1 states are 2.8, 53, 53, and 66 TeV. The lightest state, which is one order of magnitude lighter than the next to lightest one, couples to fermions through the  $\lambda^8$  flavor generator and with equal strength to left/right up/down type fermions (the unequal intensity of shading in Fig. 3 is compensated by the different values of the gauge couplings). Although its coupling

to the third generation is the largest, the lightest vector couples also to the first two generations, which makes it accessible at the LHC. For all practical purposes it corresponds to a flavor non-universal leptophobic  $Z'$ . The existing mass bounds on analogous resonances from Tevatron searches in the  $t\bar{t}$  channel lie below 1 TeV [27].

The masses of the three lightest fermion fields are 0.40, 1.8, and 90 TeV. In this case both lightest states, the  $t'$  and the  $b'$ , should be within the reach of the LHC. It is important to keep in mind however that, contrary to 4th generation quark fields, their couplings to the SM  $W$  and  $Z$  bosons arise through mixing with SM left-handed fields and are suppressed by the small angles  $s_{u_{L3}}$  and  $s_{d_{L3}}$ , which for the current choice of parameters are 0.05 and 0.02, respectively.

It is interesting to check how much this model is actually safe against existing bounds. The values of parameters chosen for this case correspond to the point with symbol “O” in the  $(\lambda_{u/d}, M_{u/d})$  planes of Figs. 1 and 2. The contributions to  $Z \rightarrow b\bar{b}$ , EW precision observables and  $V_{tb}$  read

$$\begin{aligned}\frac{\delta R_b}{R_b} &= -1.0 \cdot 10^{-3}, \\ S &= 0.00, \quad T = 0.01, \quad U = 0.00, \\ V_{tb} &= 1.00.\end{aligned}\tag{5.2}$$

Except for the correction to  $R_b$  which is naturally suppressed by the  $b$  mass the small corrections to the observables above are due to the choice  $\lambda_u \simeq y_t$ . In this region of parameters, which arises automatically anytime  $\lambda'_u \hat{Y}_u^3 \ll M_u$ , the new physics in the up-sector decouples from the SM.

As discussed in section 2.1 the exchange of flavor gauge bosons can also produce flavor breaking 4-fermion operators at tree level. Existing strong bounds on these operators are often used to rule out the possibility of low scale flavor vector fields. However the inverted hierarchy present in our spectrum allows to easily avoid all such bounds. Indeed, from Fig. 3 we note that the vector fields mediating transitions among the first and the higher generations are among the heaviest, followed by those mediating transition among the second and the third generations, while the lightest is flavor diagonal. At tree level the strongest bounds come from  $\Delta F = 2$  quark transitions, whose bounds on 4-fermion operators are conveniently summarized in [30]. Our vector boson only produce three types of such operators at tree level:

$$\begin{aligned}Q_1^{q_i q_j} &= \bar{q}_{jL}^\alpha \gamma_\mu q_{iL}^\alpha \bar{q}_{jL}^\beta \gamma^\mu q_{iL}^\beta, \\ \tilde{Q}_1^{q_i q_j} &= \bar{q}_{jR}^\alpha \gamma_\mu q_{iR}^\alpha \bar{q}_{jR}^\beta \gamma^\mu q_{iR}^\beta, \\ Q_5^{q_i q_j} &= \bar{q}_{jR}^\alpha q_{iL}^\beta \bar{q}_{jL}^\beta q_{iR}^\alpha.\end{aligned}\tag{5.3}$$

The coefficients of these operators can be obtained numerically from eq. (2.9). In our explicit example they read:

	Re (in $\text{GeV}^{-2}$ )	Im (in $\text{GeV}^{-2}$ )
$C_K^1$	$-1 \cdot 10^{-14}$	$-1 \cdot 10^{-19}$
$\tilde{C}_K^1$	$-2 \cdot 10^{-16}$	$-2 \cdot 10^{-21}$
$C_K^5$	$-5 \cdot 10^{-15}$	$-6 \cdot 10^{-20}$
$C_D^1$	$-2 \cdot 10^{-20}$	$-2 \cdot 10^{-23}$
$\tilde{C}_D^1$	$-2 \cdot 10^{-25}$	$-2 \cdot 10^{-28}$
$C_D^5$	$-2 \cdot 10^{-22}$	$-2 \cdot 10^{-25}$
$C_{B_d}^1$	$1 \cdot 10^{-16}$	$5 \cdot 10^{-16}$
$\tilde{C}_{B_d}^1$	$9 \cdot 10^{-22}$	$3 \cdot 10^{-21}$
$C_{B_d}^5$	$1 \cdot 10^{-18}$	$5 \cdot 10^{-18}$
$C_{B_s}^1$	$3 \cdot 10^{-13}$	$-4 \cdot 10^{-13}$
$\tilde{C}_{B_s}^1$	$4 \cdot 10^{-16}$	$-6 \cdot 10^{-16}$
$C_{B_s}^5$	$4 \cdot 10^{-14}$	$-6 \cdot 10^{-14}$

We have used here the notation for coefficients of Ref. [30]. Comparing with the bounds in that work,

	Re (in $\text{GeV}^{-2}$ )	Im (in $\text{GeV}^{-2}$ )
$C_K^1$	$[-9.6, 9.6] \cdot 10^{-13}$	$[-4.4, 2.8] \cdot 10^{-15}$
$\tilde{C}_K^1$	$[-9.6, 9.6] \cdot 10^{-13}$	$[-4.4, 2.8] \cdot 10^{-15}$
$C_K^5$	$[-1.0, 1.0] \cdot 10^{-14}$	$[-5.2, 2.9] \cdot 10^{-17}$
$ C_D^1 $	$< 7.2 \cdot 10^{-14}$	
$ \tilde{C}_D^1 $	$< 7.2 \cdot 10^{-14}$	
$ C_D^5 $	$< 4.8 \cdot 10^{-13}$	
$ C_{B_d}^1 $	$< 2.3 \cdot 10^{-11}$	
$ \tilde{C}_{B_d}^1 $	$< 2.3 \cdot 10^{-11}$	
$ C_{B_d}^5 $	$< 6.0 \cdot 10^{-13}$	
$ C_{B_s}^1 $	$< 1.1 \cdot 10^{-9}$	
$ \tilde{C}_{B_s}^1 $	$< 1.1 \cdot 10^{-9}$	
$ C_{B_s}^5 $	$< 4.5 \cdot 10^{-11}$	

one realizes that the resulting FCNC processes are well within the experimental bounds, with the most dangerous one ( $\text{Re } C_K^5$ ) still a factor of two smaller than current limits. This is so even if we chose extreme parameters that make flavorful new physics lie just beyond the exclusion bounds from direct searches and from flavor non-violating observables.

## 5.2 Second example: An $SU(3)^3 \times U(1)^2$ model

Our second example involves the gauging of  $SU(3)^3 \times U(1)^2$ . With respect to the previous one, two extra vector fields have been added, corresponding to the right-handed up and down flavor numbers. The flavor gauge fields can thus be identified with the generators of the  $SU(3)_{Q_L} \times U(3)_{U_R} \times U(3)_{D_R}$  group. This means that now they are free to mix (even above the flavor breaking scale) with the

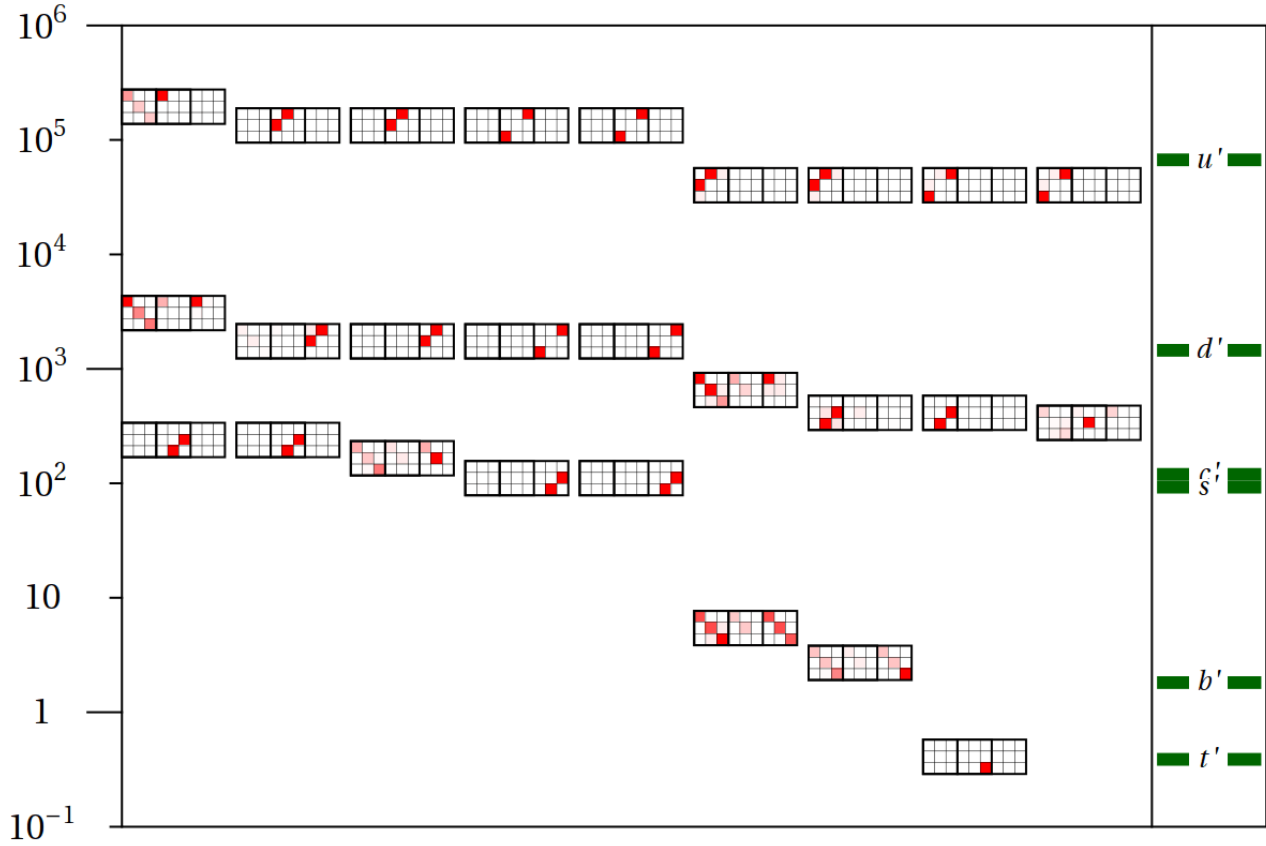


Figure 4: *Spectrum of the flavor spin-1 (left) and spin-1/2 (right) fields for the second example (see text and caption of Fig. 3 for details).*

SM hypercharge. This effect produces a mixing between the flavor gauge bosons (only those with a non-vanishing U(1) component) and the SM  $Z$  boson. We thus have two extra free parameters, characterizing the hypercharge mixing with each of the two flavor U(1). There are two expected sizes for such mixing: a)  $\mathcal{O}(1)$  if they started  $\mathcal{O}(1)$  at some high scale; b) “one-loop suppressed”  $\times$  “logs” if they were suppressed for some reason at the high scale and are produced only via radiative corrections. Since the mixing with the hypercharge does not change the flavor-breaking structure, it will not alter significantly the calculation of the flavor breaking effects. However the mixing with the hypercharge now allows such vector fields to couple to leptons at tree level. On the one hand the mixing makes it easier to detect such vector fields through their leptonic channels. On the other, however, it makes the bound on their masses stronger, in order to escape limits from electroweak precision tests.

The values of parameters we choose in this example are similar to those of the previous one:

$M_u$ (GeV)	$M_d$ (GeV)	$\lambda_u$	$\lambda'_u$	$\lambda_d$	$\lambda'_d$	$g_Q$	$g_U$	$g_D$
350	100	1.1	0.5	0.25	0.25	0.3	1	0.3

As before together with the values of the SM Yukawa couplings these parameters fix the values of the

flavon VEVs:

$$\begin{aligned} Y_u &\approx \text{Diag} (1 \cdot 10^5, 2 \cdot 10^2, 3 \cdot 10^{-1}) \cdot V \text{ TeV}, \\ Y_d &\approx \text{Diag} (6 \cdot 10^3, 4 \cdot 10^2, 7) \text{ TeV}. \end{aligned} \tag{5.4}$$

The spectrum is summarized in Fig. 4. As expected the fermionic spectrum is not very different from the previous example, with the three lightest states having masses of 0.4, 1.8, and 90 TeV.

The gauge boson spectrum underwent larger modifications. Now we have two extra states, which populate the lowest part of the spectrum. The four lightest states have now the following masses: 0.29, 1.9, 3.9, and 80 TeV. Thus we have three flavor gauge bosons that in principle are within the reach of the LHC. In particular the possibility of having vector fields associated to non-traceless generators allowed the presence of a very light vector particle coupled only to the right-handed third generation charge  $+2/3$  quarks,  $t_R$  and  $t'_R$ , since it receives a mass only from the  $\hat{Y}_u^3$  entry of the flavon field, which is the smallest one.

Neglecting possible kinetic mixing, the lightest vector couples only to  $t_R$ . For this reason it could have escaped detection and indirect bounds, despite its low mass. Once the mixing with hypercharge is taken into account strong bounds from EWPT may start becoming important: for a  $Z'$  coupling to hypercharge the bound reads:  $M_{Z'}/g_{Z'} \leq 8.55 \text{ TeV}$  (95% CL) [28]. This implies that in this explicit model we may allow for a mixing not bigger than 5% to avoid conflicts with EWPT.

The position of the parameters chosen for this example in the  $(M_{u/d}, \lambda_{u/d})$  plane is shown with the symbol “ $\times$ ” in Figs. 1 and 2, thus within the experimental bounds coming from direct searches of  $m_{b'}$  and  $m_{t'}$  and indirect effects such as  $Z \rightarrow b\bar{b}$ , EWPT and  $V_{tb}$ . In particular, for these latter quantities we get

$$\begin{aligned} \frac{\delta R_b}{R_b} &= -1.0 \cdot 10^{-3}, \\ S &= 0.00, \quad T = 0.15, \quad U = 0.01, \\ V_{tb} &= 0.97. \end{aligned} \tag{5.5}$$

This is close to the bounds from EWPT, as can also be seen from Fig. 2. The figure also shows that for this values of parameters a heavy Higgs with mass up to  $\sim 350 \text{ GeV}$  is still allowed.

Finally the effects on  $\Delta F = 2$  processes are (see the previous example for details):

	Re (in $\text{GeV}^{-2}$ )	Im (in $\text{GeV}^{-2}$ )
$C_K^1$	$-7 \cdot 10^{-15}$	$-8 \cdot 10^{-20}$
$\tilde{C}_K^1$	$-1 \cdot 10^{-16}$	$-1 \cdot 10^{-21}$
$C_K^5$	$-4 \cdot 10^{-15}$	$-4 \cdot 10^{-20}$
$C_D^1$	$-3 \cdot 10^{-20}$	$-3 \cdot 10^{-23}$
$\tilde{C}_D^1$	$-3 \cdot 10^{-25}$	$-4 \cdot 10^{-28}$
$C_D^5$	$-4 \cdot 10^{-22}$	$-4 \cdot 10^{-25}$
$C_{B_d}^1$	$2 \cdot 10^{-16}$	$2 \cdot 10^{-16}$
$\tilde{C}_{B_d}^1$	$1 \cdot 10^{-21}$	$1 \cdot 10^{-21}$
$C_{B_d}^5$	$2 \cdot 10^{-18}$	$2 \cdot 10^{-18}$
$C_{B_s}^1$	$3 \cdot 10^{-13}$	$-4 \cdot 10^{-13}$
$\tilde{C}_{B_s}^1$	$5 \cdot 10^{-16}$	$-6 \cdot 10^{-16}$
$C_{B_s}^5$	$5 \cdot 10^{-14}$	$-6 \cdot 10^{-14}$

which are similar or smaller to those of the previous example, thus well within the experimental bounds.

## 6 Discussion

We have investigated the possibility of gauging the SM flavor symmetries. Remarkably cancellation of gauge anomalies automatically leads to a model characterized by a hierarchical structure of new physics where the light generations are protected from large corrections with respect to the SM predictions, while deviations could be present for the top and bottom quarks. Contrary to the standard lore, the mechanism described here allows the scale of flavor physics to be as low as a TeV while avoiding all flavor and precision electroweak bounds but within reach at the Tevatron and the LHC. The lightest new states are the top partners in the fermionic sector and a few flavor gauge bosons that behave as non-universal  $Z'$ . Depending on the flavor gauge group a few flavor gauge bosons could be observable. Most of the spectrum is much heavier than a TeV and can not be accessed directly in present day accelerators. However the contributions could still be important for precision observables particularly in flavor physics. The actual details of the model can vary substantially (the choice of the gauge group, the number of flavon fields, values of coupling constants, etc.) however the general structure of inverted hierarchy is rather robust.

The main drawback of our model is that the scale of new physics, roughly set by the parameters  $M_{u,d}$ , is an arbitrary parameter which, if larger than few TeV would render the new states heavy, out of reach of present experiments. We are tempted to speculate that the scale of flavor physics is linked to the electroweak scale implementing this mechanism within a theory that addresses the hierarchy problem in the SM. In general flavor physics imposes formidable constraints on physics beyond the SM. At present two strategies seem possible. The first is to demand that new physics respects a MFV structure. To our knowledge however this hypothesis cannot be derived from a symmetry of the UV



theory but only arise in the IR accidentally. The other possibility is the idea of partial compositeness, see for example [31]. In this case the light generations are elementary as in the SM and the flavor transitions are suppressed by the small mixing with composite states to which the Higgs couple. Unfortunately the flavor and CP protection achieved in this case seems at present incomplete. The inverted hierarchy of our model has some similarities with both scenarios but here the suppression is due to the large mass of the relevant degrees of freedom rather than the coupling. Of course some obvious challenges should be faced in particular how to avoid reintroducing quadratic divergences in the Higgs sector once the new physics has hierarchical scales.

We hope to return to this question in the future.

**Acknowledgments:** We are grateful to Gia Dvali for inspiring discussions. We also thank Géraldine Servant for discussions, Gino Isidori and Graham Ross for comments on the manuscript and Diego Guadagnoli for pointing out some typos in appendix B. The work of BG was supported in part by the US Department of Energy under contract DE-FG03-97ER40546.

## A Radial modes

Radial and GB modes of the flavon fields can be parametrized as (see also [32] for related discussion),

$$\begin{aligned} Y_u &= U_U \rho_u U_Q^\dagger, \\ Y_d &= U_D \rho_d U_Q^\dagger, \end{aligned} \tag{A.1}$$

where  $U_{Q,U,D}$  are the three unitary matrices parametrizing the  $9+9+8=26$  Goldstone modes.  $\rho_u$  and  $\rho_d$  are the matrices of the remaining  $36-26=10$  radial modes, with VEVs

$$\begin{aligned} \langle \rho_u \rangle &= \frac{\lambda_u M_u}{\lambda'_u} \hat{g}_u^{-1} V, \\ \langle \rho_d \rangle &= \frac{\lambda_d M_d}{\lambda'_d} \hat{g}_d^{-1}, \end{aligned} \tag{A.2}$$

where, for simplicity, we assumed small Yukawa couplings (for the third generation the exact expression, Eq. (3.3), should be used). Requiring that the radial modes in  $\rho_{u,d}$  be orthogonal to the Goldstone modes correspond in the unitary gauge to the condition that cross product terms of the type  $\partial_\mu \rho A^\mu$  vanish. This correspond to the the three sets of conditions

$$\begin{aligned} A_U : \quad & \text{Im Tr}[\rho_u \partial \rho_u^\dagger \lambda^\alpha] = 0, & \alpha = 1, \dots, 9 \\ A_D : \quad & \text{Im Tr}[\rho_d \partial \rho_d^\dagger \lambda^\alpha] = 0, & \alpha = 1, \dots, 9 \\ A_Q : \quad & \text{Im Tr}[(\partial \rho_u^\dagger \rho_u + \partial \rho_d^\dagger \rho_d) \lambda^\alpha] = 0, & \alpha = 1, \dots, 8. \end{aligned} \tag{A.3}$$

We can conveniently rewrite

$$\begin{aligned} \rho_u &= \Sigma_{Ru} D_u \Sigma_{Lu}^\dagger V, \\ \rho_d &= \Sigma_{Rd} D_d \Sigma_{Ld}^\dagger, \end{aligned} \tag{A.4}$$

where  $D_{u,d}$  are real diagonal matrices parametrizing the 6 radial modes associated to the Yukawa masses ( $\langle D_{u,d} \rangle = \lambda_{u,d} M_{u,d} \hat{y}_{u,d}^{-1} / \lambda'_{u,d}$ ), and  $\Sigma_{Ru,Lu,Rd,Ld}$  are four unitary matrices parametrizing (after imposing the constraints) the remaining four angle modes. In particular we can write

$$\Sigma_X = \exp(i\Pi_X) = \exp\left(i\frac{\lambda^\alpha}{2}\pi_X^\alpha\right). \quad (\text{A.5})$$

Out of the corresponding 36 fields  $\pi_X^\alpha$  only 4 combinations remain since 6 cancel out in the combination (A.4) and 26 are removed by the constraints (A.3). The latter in terms of the  $\Pi_X$  fields read

$$\begin{aligned} A_U : \quad & \text{Tr}[(2D_u\Pi_{Lu}D_u - D_u^2\Pi_{Ru} - \Pi_{Ru}D_u^2)\lambda^\alpha] = 0, & \alpha = 1, \dots, 9 \\ A_D : \quad & \text{Tr}[(2D_d\Pi_{Ld}D_d - D_d^2\Pi_{Rd} - \Pi_{Rd}D_d^2)\lambda^\alpha] = 0, & \alpha = 1, \dots, 9 \\ A_Q : \quad & \text{Tr}[V^\dagger(2D_u\Pi_{Ru}D_u - D_u^2\Pi_{Lu} - \Pi_{Lu}D_u^2)V \\ & + 2D_d\Pi_{Rd}D_d - D_d^2\Pi_{Ld} - \Pi_{Ld}D_d^2)\lambda^\alpha] = 0, & \alpha = 1, \dots, 8. \end{aligned} \quad (\text{A.6})$$

The combination of  $\pi_X^\alpha$  which cancel out can be found from the relations

$$\Sigma_{Ru}D_u\Sigma_{Lu}^\dagger = D_u, \quad \Sigma_{Rd}D_d\Sigma_{Ld}^\dagger = D_d, \quad (\text{A.7})$$

which give the following constraints for the remaining modes:

$$\pi_{Ru}^{3,8,9} = -\pi_{Lu}^{3,8,9}, \quad \pi_{Rd}^{3,8,9} = -\pi_{Ld}^{3,8,9}. \quad (\text{A.8})$$

The first 9+9 conditions of (A.3) give the following constraints:

$$\begin{aligned} 2d_ud_c\pi_{Lu}^{1,2} &= \pi_{Ru}^{1,2}(d_u^2 + d_c^2), \\ 2d_ud_t\pi_{Lu}^{4,5} &= \pi_{Ru}^{4,5}(d_u^2 + d_t^2), \\ 2d_cd_t\pi_{Lu}^{6,7} &= \pi_{Ru}^{6,7}(d_c^2 + d_t^2), \\ \pi_{Ru}^{3,8,9} &= \pi_{Lu}^{3,8,9}, \\ 2d_d d_s \pi_{Ld}^{1,2} &= \pi_{Rd}^{1,2}(d_d^2 + d_s^2), \\ 2d_d d_b \pi_{Ld}^{4,5} &= \pi_{Rd}^{4,5}(d_d^2 + d_b^2), \\ 2d_d d_b \pi_{Ld}^{6,7} &= \pi_{Rd}^{6,7}(d_s^2 + d_b^2), \\ \pi_{Rd}^{3,8,9} &= \pi_{Ld}^{3,8,9}, \end{aligned} \quad (\text{A.9})$$

where  $\langle D_{u,d} \rangle = \text{Diag}(d_{u,d}, d_{c,s}, d_{t,b})$  and together with the previous condition imply that  $\pi_X^{3,8,9} = 0$ .

We thus ended up with 12 fields, without lack of generality  $\pi_{Ru,Rd}^{1,2,4,5,6,7}$ . There are 8 further constraints from the last line in (A.6), which leave only four independent fields. Notice that these are the only constraints that make the CKM angles appear. The expressions we obtain are quite lengthy and we do not report them here explicitly, but we only notice that all the 12 fields are in general different from zero and can be written as linear combination of four independent fields.

## A.1 General facts about radial modes

From the parametrization given above we notice some interesting facts about the way the radial modes couple. Among the invariants that can be written in the Lagrangian those which are only functions of one type of flavon field, depend only on the diagonal modes in a simple way. Indeed

$$\begin{aligned}\text{Tr}[(Y_{u,d}^\dagger Y_{u,d})^n] &= \text{Tr}[(D_{u,d}^{2n}], \\ \text{Det}[Y_{u,d}] &= \text{Det}[D_{u,d}].\end{aligned}\tag{A.10}$$

$\Sigma$ -flavons only appear when both  $Y_u$  and  $Y_d$  are present. We have for example

$$\text{Tr}[Y_u^\dagger Y_u Y_d^\dagger Y_d] = \text{Tr}[D_u^2 (\Sigma_{Lu} V \Sigma_{Ld}^\dagger) D_d^2 (\Sigma_{Lu} V \Sigma_{Ld}^\dagger)^\dagger],\tag{A.11}$$

and in general the operators will be strings of the type

$$\text{Tr}[D_u^{2n_1} (\Sigma_{Lu} V \Sigma_{Ld}^\dagger) D_d^{2n_2} (\Sigma_{Lu} V \Sigma_{Ld}^\dagger)^\dagger D_u^{2n_3} (\Sigma_{Lu} V \Sigma_{Ld}^\dagger)^\dagger \dots].\tag{A.12}$$

Thus the CKM radial modes only enter through the combination  $(\Sigma_{Lu} V \Sigma_{Ld}^\dagger)$ .

In the coupling to the SM fermions we have instead

$$\begin{aligned}\bar{Q} \tilde{H} \frac{\lambda_u M_u}{\lambda'_u} Y_u^{-1} U_R &= \bar{Q} \tilde{H} \frac{\lambda_u M_u}{\lambda'_u} V^\dagger \Sigma_{Lu} D_u^{-1} \Sigma_{Ru}^\dagger U_R, \\ \bar{Q} H \frac{\lambda_d M_d}{\lambda'_d} Y_d^{-1} D_R &= \bar{Q} H \frac{\lambda_d M_d}{\lambda'_d} \Sigma_{Ld} D_d^{-1} \Sigma_{Rd}^\dagger D_R.\end{aligned}\tag{A.13}$$

Calculating only three particle vertices, relevant for tree-level flavor breaking, we have two possible types of operators, from the interaction of SM fermions to the diagonal and the CKM radial modes respectively. For the first we can put the Higgs and the CKM modes to their VEVs  $\Sigma_X = 1$  and get

$$\begin{aligned}\frac{v}{\sqrt{2}} \bar{U}_L V^\dagger \frac{\lambda_u M_u}{\lambda'_u} D_u^{-1} U_R &\rightarrow \frac{v}{\sqrt{2}} \bar{U}_L \frac{\lambda_u M_u}{\lambda'_u} D_u^{-1} U_R = -\frac{\sqrt{2} \lambda'_u}{\lambda_u} \frac{m_{u_i}^2}{M_u v} \bar{U}_L^i D_u^{ii} U_R^i, \\ \frac{v}{\sqrt{2}} \bar{D}_L \frac{\lambda_d M_d}{\lambda'_d} D_d^{-1} D_R &= -\frac{\sqrt{2} \lambda'_d}{\lambda_d} \frac{m_{d_i}^2}{M_d v} \bar{D}_L^i D_d^{ii} D_R^i,\end{aligned}\tag{A.14}$$

where we went from the Yukawa to the quark mass eigenstate basis  $U_L \rightarrow V U_L$ . We make two observations here. First, the interactions of these modes are doubly suppressed by the Yukawa coupling constants (one suppression more than for the Higgs), and second, the interactions are flavor diagonal in the mass eigenstate basis, a sort of GIM mechanism is at work and they do not induce FC processes at tree level.

The interactions of the CKM modes read instead

$$\begin{aligned}\bar{U}_L^i \Sigma_{Lu}^{ij} m_{u_j} (\Sigma_{Ru}^\dagger)^{jk} U_R^k &= i \bar{U}_L^i (\Pi_{Lu}^{ij} m_{u_j} - m_{u_i} \Pi_{Ru}^{ij}) U_R^j, \\ \bar{D}_L^i \Sigma_{Ld}^{ij} m_{d_j} (\Sigma_{Rd}^\dagger)^{jk} D_R^k &= i \bar{D}_L^i (\Pi_{Ld}^{ij} m_{d_j} - m_{d_i} \Pi_{Rd}^{ij}) D_R^j,\end{aligned}\tag{A.15}$$

from where we see that they can mediate flavor violations. Alternatively, the  $\Sigma$ -fields can be reabsorbed into a field redefinition of the quark fields, which makes the interactions appear from the kinetic terms:

$$\begin{aligned} i\bar{U}_{L,R}\gamma^\mu\Sigma_{Lu,Ru}^\dagger\partial_\mu\Sigma_{Lu,Ru}U_{L,R} &= -\bar{U}_{L,R}\gamma^\mu\partial_\mu\Pi_{Lu,Ru}U_{L,R}, \\ i\bar{D}_{L,R}\gamma^\mu\Sigma_{Ld,Rd}^\dagger\partial_\mu\Sigma_{Ld,Rd}D_{L,R} &= -\bar{D}_{L,R}\gamma^\mu\partial_\mu\Pi_{Ld,Rd}D_{L,R}. \end{aligned} \quad (\text{A.16})$$

In this form the interactions of the CKM modes resembles the one of the longitudinal modes of the vector fields. To estimate the potential flavor violation induced by these interactions, we should write explicitly the dependence on the independent modes in the  $\Pi$ -fields and work out the spectrum of the corresponding modes. The full analytic expression turns out to be lengthy and not very illuminating, therefore we will give them explicitly in the two-flavor case to illustrate their structure, while for the three-flavor case we will give only the numerical estimates.

## A.2 The 2-flavors example

In the two dimensional case there is only one CKM mode, which means that all  $\Pi$ -fields can be rewritten in terms of one field only. The diagonal entries vanish because of the constraints, like in the 3 flavor case. It is simple to work out all the constraints and the explicit results for the  $\Pi$ -fields, in terms of the canonically normalized field  $\varphi$ , read

$$\begin{aligned} \Pi_{Lu} &= \frac{\sigma^2}{2} \frac{d_d^2 - d_s^2}{d_u^2 - d_c^2} \frac{d_u^2 + d_c^2}{\kappa} \varphi, \\ \Pi_{Ru} &= \frac{\sigma^2}{2} \frac{d_d^2 - d_s^2}{d_u^2 - d_c^2} \frac{2d_u d_c}{\kappa} \varphi, \\ \Pi_{Ld} &= \frac{\sigma^2}{2} \frac{d_u^2 - d_c^2}{d_d^2 - d_s^2} \frac{d_d^2 + d_s^2}{\kappa} \varphi, \\ \Pi_{Rd} &= \frac{\sigma^2}{2} \frac{d_u^2 - d_c^2}{d_d^2 - d_s^2} \frac{2d_d d_s}{\kappa} \varphi, \\ \kappa &= \sqrt{(d_d^2 + d_s^2)(d_u^2 - d_c^2)^2 + (d_u^2 + d_c^2)(d_d^2 - d_s^2)^2}, \end{aligned} \quad (\text{A.17})$$

where  $\sigma^a$  are Pauli matrices. As we explained earlier only the combination  $(\Sigma_{Lu}V\Sigma_{Ld}^\dagger)$  can appear in the scalar potential. Notice that in this case  $V = \exp(i\sigma^2\theta_{12})$  and therefore

$$(\Sigma_{Lu}V\Sigma_{Ld}^\dagger) = \exp(i(\Pi_{Lu} - \Pi_{Ld} + \sigma^2\theta_{12})) = \exp(i\sigma_{12}(\theta_{12} + \frac{\kappa}{2(d_u^2 - d_c^2)(d_d^2 - d_s^2)}\varphi)), \quad (\text{A.18})$$

so that the CKM modes in this case enter like a shift of the Cabibbo angle in the scalar potential.

The interactions with the SM fermions read instead

$$\begin{aligned} \frac{1}{2}\bar{u}_L \frac{d_d^2 - d_s^2}{d_u^2 - d_c^2} \frac{d_u^2 + d_c^2}{\kappa} m_c (1 - \frac{2d_c^2}{d_u^2 + d_c^2}) c_R \varphi &\simeq \frac{1}{2} \frac{d_d}{d_u \sqrt{d_u^2 + d_d^2}} m_c \bar{u}_L c_R \varphi \approx \frac{\sqrt{2}\lambda'_u m_u m_c}{\lambda_u M_u v} \bar{u}_L c_R \varphi \\ \frac{1}{2}\bar{c}_L \frac{d_d^2 - d_s^2}{d_u^2 - d_c^2} \frac{d_u^2 + d_c^2}{\kappa} m_u (\frac{2d_u^2}{d_u^2 + d_c^2} - 1) u_R \varphi &\simeq \frac{1}{2} \frac{d_d}{d_u \sqrt{d_u^2 + d_d^2}} m_u \bar{c}_L u_R \varphi \approx \frac{\sqrt{2}\lambda'_u m_u^2}{\lambda_u M_u v} \bar{c}_L u_R \varphi \end{aligned} \quad (\text{A.19})$$

and analogously for the down-type quarks. The most dangerous interaction is suppressed by  $\frac{\lambda'_u m_u m_c}{\lambda_u M_u v}$  which, like for the diagonal modes, provide an extra Yukawa suppression with respect to the Higgs coupling, which is already Yukawa suppressed.

In this case the smallness of the coupling guarantees no dangerous tree-level FC effects regardless of what the masses of the radial modes may be, as long as they are above the bounds from direct searches (and such bounds can be even quite loose because of the small couplings).

### A.3 The 3-flavors case: numerical

In the 3-flavor case the formulae are lengthy and less intuitive, however one can still calculate numerically the 3-fields vertices involving two SM fermions and one of the four CKM radial modes (canonically normalized). Even without knowing the potential, and therefore the mass matrix of these radial modes, one can estimate their maximum FC contributions by assuming that the lightest eigenmode couples to the vertices with the largest couplings. In this case contributions to  $\Delta F = 2$  operators of the form

$$\frac{c}{m_\pi^2} (\bar{q}q)^2$$

are obtained, with

	Re( $c$ )	Im( $c$ )
$(\bar{s}_R d_L)^2$	$-7 \cdot 10^{-18}$	$-1 \cdot 10^{-20}$
$(\bar{c}_R u_L)^2$	$-4 \cdot 10^{-18}$	$-3 \cdot 10^{-19}$
$(\bar{b}_R d_L)^2$	$-7 \cdot 10^{-17}$	$-7 \cdot 10^{-22}$
$(\bar{b}_R s_L)^2$	$-8 \cdot 10^{-13}$	$-9 \cdot 10^{-18}$

and with  $m_\pi$  the mass of the lightest CKM mode eigenstate. The contributions are so suppressed that  $m_\pi$  can be as light as 100 MeV without incurring into problems with flavor. The quantitative results above nicely fit with what is observed in the two flavor case, and the couplings of the canonically normalized CKM modes to the fermions are numerically compatible with the short-hand formula

$$\pi_{\text{CKM}} \frac{\lambda'_{u,d} m_{q_i} m_{q_j}}{\lambda_{u,d} v M_{u,d}} \bar{q}_R^i q_L^j, \quad (\text{A.20})$$

which is similar to the flavor preserving one for the radial modes of the diagonal flavon fields.

## B Oblique Corrections

In this appendix we derive the one loop expression for the  $S$ ,  $T$  and  $U$  parameters in our model.

We use the standard definitions [33],

$$\begin{aligned} S &= -16\pi \Pi'_{3Y}(0), \\ T &= \frac{4\pi}{s_w^2 c_w^2 M_Z^2} [\Pi_{11}(0) - \Pi_{33}(0)], \\ U &= 16\pi [\Pi'_{11}(0) - \Pi'_{33}(0)]. \end{aligned} \quad (\text{B.1})$$

For simplicity we work in the limit where only the mixing of the top is important. From the couplings (3.8), (3.10) the contribution of the third generation to  $T$  (obtained as the difference between the correlators in our model and their SM values corresponding to  $s_{u_{L3}} = 0$ ) is given by,

$$T = \frac{3\pi}{s_w^2 c_w^2 M_Z^2} [2s_{u_{L3}}^2 \Pi_{LL}(m_{t'}, m_b, 0) - 2s_{u_{L3}}^2 \Pi_{LL}(m_t, m_b, 0) + (1 - c_{u_{L3}}^4) \Pi_{LL}(m_t, m_t, 0) - s_{u_{L3}}^4 \Pi_{LL}(m_{t'}, m_{t'}, 0) - 2s_{u_{L3}}^2 c_{u_{L3}}^2 \Pi_{LL}(m_{t'}, m_t, 0)] \quad (\text{B.2})$$

where we have introduced the self energies  $\Pi_{LL}(m_1, m_2, q)$  with two left currents. In the limit  $m_b \rightarrow 0$  one finds,

$$T = \frac{3s_{u_{L3}}^2}{8\pi s_w^2 c_w^2} \frac{m_t^2}{M_Z^2} \left[ c_{u_{L3}}^2 \left( \frac{m_{t'}^2}{m_{t'}^2 - m_t^2} \log\left(\frac{m_{t'}^2}{m_t^2}\right) - 1 \right) + \frac{s_{u_{L3}}^2}{2} \left( \frac{m_{t'}^2}{m_t^2} - 1 \right) \right]. \quad (\text{B.3})$$

Repeating the same steps for  $S$  one obtains,

$$S = 4\pi [s_{u_{L3}}^2 (3s_{u_{L3}}^2 - 2) \Pi'_{LL}(m_t, m_t, 0) + s_{u_{L3}}^2 (3s_{u_{L3}}^2 - 4) \Pi'_{LL}(m_{t'}, m_{t'}, 0) + 6s_{u_{L3}}^2 c_{u_{L3}}^2 \Pi'_{LL}(m_t, m_{t'}, 0) + 4s_{u_{L3}}^2 \Pi'_{LR}(m_t, m_t, 0) - 4s_{u_{L3}}^2 \Pi'_{LR}(m_{t'}, m_{t'}, 0)], \quad (\text{B.4})$$

which gives

$$S = \frac{s_{u_{L3}}^2}{6\pi} \left[ \left( 3c_{u_{L3}}^2 \frac{(m_{t'}^2 + m_t^2)(m_{t'}^4 - 4m_{t'}^2 m_t^2 + m_t^4)}{(m_{t'}^2 - m_t^2)^3} - 1 \right) \log\left(\frac{m_{t'}^2}{m_t^2}\right) - c_{u_{L3}}^2 \frac{5m_{t'}^4 - 22m_{t'}^2 m_t^2 + 5m_t^4}{(m_{t'}^2 - m_t^2)^2} \right]. \quad (\text{B.5})$$

For completeness the  $U$  parameter is given by,

$$U = \frac{s_{u_{L3}}^2}{6\pi} \left[ -3 \left( c_{u_{L3}}^2 \frac{(m_{t'}^2 + m_t^2)(m_{t'}^4 - 4m_{t'}^2 m_t^2 + m_t^4)}{(m_{t'}^2 - m_t^2)^3} - 1 \right) \log\left(\frac{m_{t'}^2}{m_t^2}\right) + c_{u_{L3}}^2 \frac{5m_{t'}^4 - 22m_{t'}^2 m_t^2 + 5m_t^4}{(m_{t'}^2 - m_t^2)^2} \right]. \quad (\text{B.6})$$

For our analysis we have used the recent analysis [34],

$$\begin{aligned} S &= 0.02 \pm 0.11 \\ T &= 0.05 \pm 0.12 \\ U &= 0.07 \pm 0.12 \end{aligned} \quad \text{with correlation matrix} \quad \begin{pmatrix} 1 & 0.879 & -0.469 \\ 0.879 & 1 & -0.716 \\ -0.469 & -0.716 & 1 \end{pmatrix}. \quad (\text{B.7})$$

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